

UH-511-833-95

SUPERSYMMETRY: WHERE IT IS AND HOW TO FIND IT ^a

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We present a pedagogical, but by no means complete, review of weak scale supersymmetry phenomenology. After a general introduction to the new particles that must be present in any supersymmetric framework, we describe how to write down their interactions with one another as well as with the particles of the Standard Model. We then elucidate the assumptions underlying the Minimal Supersymmetric Model as well as the more restrictive minimal supergravity GUT model with the radiative breaking of electroweak symmetry. These models serve to guide our thinking about the implications of supersymmetry for experiments. To facilitate our study of signatures of supersymmetric particles at high energy colliders, we describe the decay patterns of sparticles as well as their production mechanisms in e^+e^- and hadron-hadron collisions. We then discuss how sparticles may be searched for in on-going experiments at the Tevatron and at LEP. We review phenomenological constraints on supersymmetric particle masses from non-observation of any signals in these experiments, and also briefly discuss constraints from low energy experiments and from cosmology. Next, we study new strategies by which supersymmetric particles may be searched for at supercolliders, and also what we can learn about their properties (masses, spins, couplings) in these experiments. A determination of sparticle properties, we will see, may provide us with clues about the nature of physics at the ultra-high scale. After a brief discussion of possible extensions of the minimal framework and the implications for phenomenology, we conclude with our outlook for the future.

^aLectures presented at the 1995 Theoretical Advanced Study Institute, University of Colorado, Boulder, USA

1 Introduction and Prelude

As Bagger has already discussed in his lectures¹, supersymmetry^{2,3,4} differs from familiar symmetries such as rotation or Lorentz invariance, gauge invariance, or the old-fashioned isospin invariance of strong interactions in one important aspect. Unlike these “bosonic” symmetries which relate properties of a boson (fermion) with those of other bosons (fermions) a supersymmetry (SUSY) inter-relates the properties of bosons and fermions, and so, provides a new level of synthesis. The fermionic generator Q of supersymmetry transforms as a spinor under the Lorentz group. It commutes with the translation generator. As a result, all states (other than the zero energy ground state) come in degenerate pairs—for each bosonic state there is a fermionic state with the same (non-zero) energy. Particularizing to single particle states, we see that all particles must have supersymmetric partners (sparticles) with the same mass but spin differing by $\frac{1}{2}$. Furthermore, because the generator of supersymmetry commutes with internal symmetry generators, sparticles must have the same gauge quantum numbers as their ordinary particle partners. Thus, aside from mixing effects, *the gauge interactions of sparticles are completely fixed*. This is the principal reason why supersymmetric models have predictive power.

Of course, this also means that unbroken supersymmetry is excluded by experiment. We know that there are no bosons with, for instance, the mass and charge of the electron. Such bosons, if they existed, would have been produced via electromagnetic interactions in accelerator experiments, and would have been long since discovered. Supersymmetry must, therefore, be a broken symmetry. In this case, it is reasonable to ask why we should bother with it at all. In this context, we must recall an important lesson we have learnt from electroweak theory: although mass relationships implied by the symmetry may be badly violated when this symmetry is spontaneously broken, the underlying relationships between couplings are nonetheless preserved, with important implications^b for physics.

The physics of supersymmetry breaking is, however, not yet understood. Nevertheless, if we believe that the main motivation¹ for *weak scale* SUSY comes from the observation that it can protect scalar masses from large radiative corrections, and thus ameliorate the fine-tuning problem of the SM, we must also accept that SUSY breaking interactions must be “soft”; *i.e.* they do not reintroduce the quadratic divergences that SUSY was introduced to elimi-

^bThis does not mean that there are physics implications only when coupling constant relationships are maintained: knowledge of how the symmetry breaking terms transform, yields observable consequences from the underlying (broken) symmetry. The isospin symmetry of strong interactions (broken by electromagnetic interactions) and flavour $SU(3)$ are, perhaps, the best known particle physics examples of this.

nate in the first place. Fortunately, the dimensionless couplings of sparticles to gauge bosons and their supersymmetric counterparts are not soft, so that the predictive power of SUSY models referred to above is unaffected by the introduction of soft SUSY-breaking terms.^c Differences between various low-energy models arise in the choice of the superpotential function¹ which gives rise to Yukawa interactions, as well as assumptions about the soft SUSY-breaking terms.

The plan of these lectures is as follows. We discuss generalities about supersymmetric particles and their interactions in Sec. 2. The assumptions underlying the Minimal Supersymmetric Model and the popular supergravity framework extensively used for phenomenology are discussed in the following two sections. Supersymmetric model-building⁵ is beyond the scope of these lectures. We will only focus on those elements that are necessary for our understanding of the phenomenology. In Sec. 5 and Sec. 6 we describe sparticle decay patterns and production mechanisms, respectively. Sec. 7 briefly outlines the available computer programs for the simulation of supersymmetric events at colliders. Empirical constraints on sparticle masses are discussed in Sec. 8. In Sec. 9 we investigate how supersymmetry might be discovered at future colliders^d while the following section focusses on how we might determine the properties (masses, spins, couplings) of sparticles. In Sec. 11 we briefly discuss possible extensions of the minimal framework, and their impact on SUSY phenomenology. Finally, we conclude with a summary and outlook for the future.

2 The Supersymmetric Framework.

2.1 Particle Content

We begin by considering the field content of the minimal supersymmetric extension of the SM. Each chiral fermion $f_{L,R}$ in the SM has a spin zero supersymmetric partner, the sfermion $\tilde{f}_{L,R}$ ($f = q, \ell$) with the same gauge quantum numbers, so that the number of bosonic and fermionic degrees of freedom are the same. Note that this means that for each massive Dirac fermion there are *two* distinct scalar partners (\tilde{f}_L and \tilde{f}_R), each represented by a complex scalar

^cThat this must be so can be simply inferred if we recall that the internal quantum numbers completely fix the gauge interactions of particles in any field theory, regardless of supersymmetry.

^dA comprehensive numerical discussion of the branching fractions, production cross sections and signal cross sections would require too many figures and tables and it is not possible to include all these here. We will, therefore, not make an attempt to display numerical results in these notes, but provide the reader with an extensive bibliography to the literature where these may be found.

field. The chiral fermion field with the corresponding scalar field and the auxiliary field necessary to linearly realize the SUSY algebra¹ together constitute a chiral matter supermultiplet, in much the same way that the proton and the neutron together constitute the isospin doublet. The complex Higgs bosons together with their spin $\frac{1}{2}$ partners, the Higgsinos (and the associated auxiliary field) comprise the Higgs chiral supermultiplets. Finally, the partner of a massless spin-1 real gauge field is a spin $\frac{1}{2}$ Majorana^e gaugino which together with a real auxiliary field constitute the gauge supermultiplet.^f

All SUSY models must, therefore, include quark and lepton chiral superfields (a superfield¹ is simply a supermultiplet whose components are fields, in exactly the same way that the Yang-Mills field is a gauge multiplet whose components are (gauge) fields) with gauge quantum numbers corresponding to those of the corresponding quark or lepton. For instance, there is an $SU(3)$ triplet, $SU(2)$ singlet superfield \hat{U}_R with $Y = \frac{4}{3}$ (the subscript R reminds us of the chirality of the corresponding SM fermion) *etc.* Since the superpotential is required to be a function of only left (or equivalently, only right) chiral superfields, we will instead of \hat{U}_R work with the anti-quark $SU(2)$ singlet superfield \hat{U}^c (henceforth, we drop the redundant chirality subscript since all fields are left-chiral), and likewise introduce weak isosinglet superfields \hat{D}^c and \hat{E}^c in addition to the quark and lepton doublets \hat{Q} and \hat{L} , respectively. As in the SM, the replication of generations must be put in by hand. We must also include gauge super-multiplets which transform according to the adjoint representation of the gauge group. Finally, we have to introduce at least two different Higgs supermultiplets \hat{h} and \hat{h}' to give masses to the up- and down- type fermions, respectively.^g We will use the notation of the Korea Lectures⁶, and require \hat{h} (\hat{h}') transform as the 2 (2^*) representation of $SU(2)$. The Minimal Supersymmetric Model (MSSM) contains the smallest number of new fields — the matter and gauge multiplets along with exactly two Higgs doublets.

2.2 Supersymmetric Interactions

Bagger has already discussed how to construct supersymmetric Lagrangians in his lectures.¹ We will not repeat this discussion here, but only recapitulate the

^eA Majorana spinor ψ is one whose charge conjugate $\psi^c = \psi$. Notice that this means that ψ_L and ψ_R^* are linearly related.

^fWe choose the Wess-Zumino gauge discussed by Bagger. Again, note the equality of the physical bosonic and fermionic degrees of freedom.

^gThe Yukawa interactions come from superpotential interactions. Since the superpotential is required to contain only left-chiral superfields, we cannot use the the Higgs field and its conjugate (which will have opposite chirality) simultaneously in the superpotential. This is unlike the situation in the SM where the sole Higgs field ϕ and its conjugate ϕ^* can be simultaneously introduced into the Lagrangian.

necessary results. The Lagrangian for global supersymmetry, after elimination of auxiliary fields, can be written in four-component spinor notation as,⁶

$$\begin{aligned}
\mathcal{L} = & \sum_i (D_\mu S_i)^\dagger (D^\mu S_i) + \frac{i}{2} \sum_i \bar{\psi}_i \not{D} \psi_i \\
& - \frac{1}{4} \sum_A F_{\mu\nu A} F_A^{\mu\nu} + \frac{i}{2} \sum_A \bar{\lambda}_A \not{D} \lambda_A \\
& - \sqrt{2} \sum_i \left[S_i^\dagger(t_A) \bar{\psi}_i \frac{1-\gamma_5}{2} \lambda_A + h.c. \right] \\
& - \frac{1}{2} \sum_A \left[\sum_i S_i^\dagger t_A S_i + \xi_A \right]^2 \\
& - \sum_i \left| \frac{\partial f}{\partial S_i} \right|^2 \\
& - \frac{1}{2} \sum_{i,j} \left\{ \bar{\psi}_i \left[\frac{1-\gamma_5}{2} \right] \frac{\partial^2 f}{\partial S_i \partial S_j} \psi_j + h.c. \right\} \tag{1}
\end{aligned}$$

Here, S_i (ψ_i) denotes the scalar (Majorana fermion^h) component of the i th chiral superfield, $F_{\mu\nu A}$ is the Yang-Mills gauge field, and λ_A is the Majorana gaugino superpartner of the corresponding gauge boson and ξ_A are constants which can be non-zero only for U(1) factors of the gauge group.⁷ In anticipation of simple grand unification, we will set these to zero. The function f in the last two lines of Eq. (1) is the superpotential which is a function of chiral superfields. By $\frac{\partial f}{\partial S_i}$ (and other derivatives), we mean differentiate with respect to \hat{S}_i and then set $\hat{S}_i = S_i$.

We note the following:

1. The first two lines are the gauge invariant kinetic energies for the components of the chiral and gauge superfields. The derivatives that appear are gauge covariant derivatives appropriate to the particular representation in which the field belongs. For example, if we are talking about SUSY QCD, for quark fields in the first line of Eq. (1) the covariant derivative contains triplet SU(3) matrices, whereas the covariant derivative acting on the gauginos in the following line will contain octet matrices. As stressed above, these terms completely determine how particles interact with gauge bosons.

^hThe four-component spinor is defined⁶ by choosing the right chiral component such that the spinor is Majorana.

2. The next line describes the interactions of gauginos with matter and Higgs multiplets. Notice that these interactions are also determined by the gauge couplings. Here t_A is the appropriate dimensional matrix representation of the group generators times the gauge coupling constant. Matrix multiplication is implied. To see that these terms are gauge invariant, recall that ψ_{iR} which is fixed by the Majorana condition, transforms according to the conjugate representation to ψ_{iL} .
3. Line four describes the quartic couplings of scalar matter. Notice that these are determined by the gauge interactions. The interactions on this line are referred to as D -terms.
4. Finally, the last two lines in Eq. (1) describe the non-gauge superpotential interactions of matter fields and lead to the Yukawa interactions responsible for matter fermion masses in the SM. Since these interactions do not involve any spacetime derivatives, choosing the superpotential to be a globally gauge invariant function of superfields is sufficient to guarantee the gauge invariance of the Lagrangian. For a renormalizable theory, the superpotential must be a polynomial of degree ≤ 3 .

Assuming the minimal field content discussed above and neglecting inter-generational mixing, the most general $SU(3) \times SU(2) \times U(1)$ invariant superpotential can be written as,

$$f = f_1 + g_1 + g_2, \quad (2)$$

with

$$\begin{aligned} f_1 = & \mu(\hat{h}^0 \hat{h}'^0 + \hat{h}^+ \hat{h}'^-) + f_u(\hat{u} \hat{h}^0 - \hat{d} \hat{h}^+) \hat{U}^c \\ & f_d(\hat{u} \hat{h}'^- + \hat{d} \hat{h}'^0) \hat{D}^c + f_e(\hat{\nu} \hat{h}'^- + \hat{e} \hat{h}'^0) \hat{E}^c + \dots, \end{aligned} \quad (3)$$

$$g_1 = \sum_{i,j,k} \left[\lambda_{ijk} \hat{L}_i \hat{L}_j \hat{E}_k^c + \lambda'_{ijk} \hat{L}_i \hat{Q}_j \hat{D}_k^c \right], \quad (4)$$

and,

$$g_2 = \sum_{i,j,k} \lambda''_{ijk} \hat{U}_i^c \hat{D}_j^c \hat{D}_k^c. \quad (5)$$

Other models may be constructed by introducing additional fields into the superpotential.

In Eq. (3), \hat{u} and \hat{d} denote the doublet quark superfields. A similar notation is used for leptons. The minus sign in the second term is because it is the anti-symmetric combination of two doublets that forms an SU(2) singlet. Since \hat{h}' is defined to transform according to the 2^* representation, the symmetric combination appears in other terms. Also, f_u , f_d and f_e are the coupling constants for the Yukawa interactions that give rise to first generation quark and lepton masses. The ellipses denote similar terms for other generations. In the Eq. (4) and (5), i, j and k denote generation indices, while the λ 's are coupling constants.^{*i*} We have, for brevity, not expanded out the gauge invariant product of doublets in Eq. (4). The Lagrangian interactions can now be obtained by substituting the superpotential (2) into Eq. (1). It is easy to check that the terms obtained from g_1 and g_2 lead to the violation of lepton and baryon number conservation, respectively. This can also be seen directly from the superpotential: for instance, with the usual assignment of lepton number of one unit to \hat{L} and \hat{E} (so that f_1 remains invariant), g_1 clearly is not globally invariant under the corresponding U(1) transformations.

This situation is quite different from the SM where the gauge invariance of the Lagrangian guarantees the absence of renormalizable baryon or lepton number violating interactions. It is the presence of scalar baryon and lepton superpartners that now allow for renormalizable baryon and lepton number violating fermion-fermion-scalar vertices. The simultaneous presence of all such terms with large couplings would lead to proton decay at the weak interaction rate if, as Bagger has explained, the superpartners have masses below the TeV scale. This would, of course, be a phenomenological disaster. Unlike as in the SM, an additional global symmetry needs to be put in by hand (or some dimensionless couplings need to be chosen to be tiny) to prevent this. One possible way^{*j*} to guarantee proton stability is to assume that at least one (or both) of baryon or lepton number is conserved (and in the case of baryon number violation, also that the lightest SUSY fermion is heavier than the proton). Within the MSSM it is assumed that both g_1 and g_2 vanish; *i.e.* the model is minimal in that it not only contains the fewest new particles, but also the fewest number of interactions necessary to be phenomenologically viable. It is easy to check that the interactions then multiplicatively conserve a new quantum number called *R*-parity which is defined to be +1 for SM particles such as quarks, leptons and gauge and Higgs bosons, and -1 for their

^{*i*}We have omitted a bilinear term in g_1 . This term can be rotated away in the supersymmetry limit.

^{*j*}It has, however, recently been argued⁸ that certain products of B- and L-violating interactions, even if simultaneously present, are only weakly constrained.

supersymmetric partners.^k It is also possible to construct phenomenologically viable models that include g_1 or g_2 terms in the superpotential and so violate R -parity conservation. We will return to such non-minimal models at the end of these lectures but will focus, for now, on the MSSM.

2.3 Supersymmetry Breaking

The interactions defined by the Lagrangian (1) are exactly supersymmetric, and so, cannot be the whole story. The physics of supersymmetry breaking is, however, not understood so that the best that we can hope for at present is a parametrization of SUSY-breaking effects. The guiding principle, as we have already noted, is that the SUSY breaking terms should not destabilize scalar masses by reintroducing the quadratic divergences that SUSY was introduced to eliminate in the first place. Girardello and Grisaru⁹ have classified all renormalizable soft SUSY breaking operators. For our purposes, it is sufficient to know that these consist of,

- explicit masses for the scalar members of chiral multiplets; *i.e.* squarks, sleptons and Higgs bosons,
- explicit masses μ_1 , μ_2 and μ_3 for the U(1), SU(2) and SU(3) gauginos,
- new super-renormalizable scalar interactions: for each trilinear (bilinear) term in the superpotential of the form $C_{ijk}\hat{S}_i\hat{S}_j\hat{S}_k$ ($C_{ij}\hat{S}_i\hat{S}_j$), we can introduce a soft supersymmetry breaking scalar interaction $A_{ijk}C_{ijk}S_iS_jS_k$ ($B_{ij}C_{ij}S_iS_j$) where the A 's and B 's are constants. These terms are often referred to as A - and B -terms.

The scalar and gaugino masses obviously serve to break the undesired degeneracy between the masses of sparticles and particles. We will see later that the explicit trilinear scalar interactions mainly affect the phenomenology of third generation sfermions.

3 The Minimal Supersymmetric Model

The MSSM is the simplest supersymmetric extension of the SM in that it contains the fewest number of fields and superpotential interactions. Here, we identify the particles, *i.e.* the mass eigenstates of the model, and also summarize the model parameters that have been introduced.

^kNotice that R -parity is automatically conserved by the interactions of gauge bosons and gauginos on the first four lines of Eq. (1). Whether or not it is a good symmetry then depends on the choice of superpotential. Spontaneous R -violation via a vacuum expectation value of a doublet sneutrino is excluded by the measurement of the Z width at LEP, as we will discuss later.

3.1 Mass Eigenstates and their Interactions

SUSY Scalars: The scalar partners \tilde{f}_L and \tilde{f}_R have the same electric charge and colour, and so can mix if $SU(2) \times U(1)$ is broken. It is simple to check that the gauge interactions conserve chiral flavour in that they couple only left (right) multiplets with one another, *i.e.* \tilde{f}_L couples only to \tilde{f}_L (f_L) via gauge boson (gaugino) interactions. Unless this “extended chiral symmetry” is broken, there can be no $\tilde{f}_L - \tilde{f}_R$ mixing. This symmetry is, however, explicitly broken by the Yukawa interactions in the superpotential.^{*l*} We thus conclude that $\tilde{f}_L - \tilde{f}_R$ mixing is proportional to the corresponding Yukawa coupling and hence to the corresponding *fermion* mass. This mixing is generally negligible except for the case of top squarks where it plays a very important role. We will, therefore, neglect this intra-generational mixing for the first five squark flavours, and for simplicity, also any inter-generational mixing.

SUSY Fermions: The gauginos and Higgsinos are the only spin- $\frac{1}{2}$ fermions. Of these, the gluinos being the only colour octet fermions, remain unmixed and have a mass $m_{\tilde{g}} = |\mu_3|$.

Electroweak gauginos and Higgsinos of the same charge can mix, once electroweak gauge invariance is broken. The mass matrices can be readily worked out using Eq. (1), (2) and (3). These matrices for both the charged and neutral electroweak gaugino sector are explicitly given elsewhere^{*6*} and will not be rewritten here. Note the translation, $\mu \equiv -2m_1$. The mass eigenstates can be obtained by diagonalizing these matrices.^{*m*} In the MSSM the charged Dirac Higgsino (composed of the charged components of the doublets \tilde{h} and \tilde{h}') and the charged gaugino (the partner of the charged W boson) mix to form two Dirac charginos, \tilde{W}_1 and \tilde{W}_2 , while the two neutral Higgsinos and the neutral $SU(2)$ and $U(1)$ gauginos mix to form four Majorana neutralinos $\tilde{Z}_1 \dots \tilde{Z}_4$, in order of increasing mass. In general, the mixing patterns are complex and depend on several parameters: μ , $\mu_{1,2}$ and $\tan \beta \equiv \frac{v}{v'}$, the ratio of the vacuum expectation values of the two Higgs fields introduced above. If either $|\mu|$ or $|\mu_1|$ and $|\mu_2|$ are very large compared to M_W , the mixing becomes small. For $|\mu| \gg M_W, |\mu_{1,2}|$, the lighter chargino is essentially a gaugino while the heavier one is a Higgsino with mass $|\mu|$; also, the two lighter neutralinos are gaugino-like while $\tilde{Z}_{3,4}$ are dominantly Higgsinos with mass $\sim |\mu|$. If instead,

^{*l*}Without the assumption of R -parity conservation there would also be mixing between the \hat{h} and \hat{L} supermultiplets. Such a mixing which is absent in the MSSM can have significant phenomenological impact.

^{*m*}If an eigenvalue of the mass matrix for any state ψ turns out to be negative, one can always define a new spinor $\psi' = \gamma_5 \psi$ which will have a positive mass. For neutralinos, ψ' should be defined with an additional factor i to preserve its Majorana nature under the γ_5 transformation.

the gaugino masses are very large, it is the heavier chargino and neutralinos that become gaugino-like.

Without further assumptions, the three gaugino masses are independent parameters. It is, however, traditional to assume that there is an underlying grand unification, and that these masses derive from a common gaugino mass parameter defined at the unification scale. The differences between the various gaugino masses then come from the fact that they have different interactions, and so, undergo different renormalization when these are evolved down from the GUT scale to the weak scale. The gaugino masses are then related by,

$$\frac{3\mu_1}{5\alpha_1} = \frac{\mu_2}{\alpha_2} = \frac{\mu_3}{\alpha_3}. \quad (6)$$

Here the α_i are the fine structure constants for the different factors of the gauge group. With this GUT assumption, \tilde{W}_1 and $\tilde{Z}_{1,2}$ will be substantially lighter than gluinos. It is for this reason that future e^+e^- colliders operating at $\sqrt{s} \simeq 500$ -1000 GeV are expected to be competitive with hadron supercolliders such as the LHC which has much higher energy. We also mention that for not too small values of $|\mu|$, the lightest neutralino tends to be dominantly the hypercharge gaugino.

The Electroweak Symmetry Breaking Sector: Although this is not in the mainstream of what we will discuss, we should mention that because there are two doublets in the MSSM, after the Higgs mechanism there are five physical spin zero Higgs sector particles left over in the spectrum. Assuming that there are no CP violating interactions in this sector, these are two neutral CP even eigenstates (H_ℓ and H_h) which behave as scalars as far as their couplings to fermions go (the subscripts ℓ and h denote light and heavy), a neutral “pseudoscalar” CP odd particle H_p , and a pair of charged particles H^\pm .

The Higgs boson sector¹⁰ of the MSSM is greatly restricted by SUSY. In addition to $\tan\beta$ which also enters the gaugino-Higgsino sector, the tree level Higgs sector is fixed by just one additional parameter which may be taken to be m_{H_p} . In particular, the Higgs quartic self-couplings are all given by those on line four of Eq. (1) and so are fixed to be $O(g^2)$. This leads to the famous (tree-level) bound, $m_{H_\ell} < \min[M_Z, m_{H_p}]|\cos 2\beta|$. This receives important corrections from t and \bar{t} loops because of the rather large value of the top Yukawa coupling and the bound is weakened¹¹ to about 120-130 GeV depending on the value of m_t . Thus, in contrast to early expectations, H_ℓ may well escape detection at LEP2. It is worth mentioning that if we assume that all couplings remain perturbative up to the GUT scale, then the mass of the lightest Higgs boson is bounded by 145-150 GeV in *any* weak scale SUSY model.¹² The physics behind this is the same as that behind the bound¹³

$m_{H_{SM}} \lesssim 200$ GeV on the mass of the SM Higgs boson, obtained under the assumption that the Higgs self-coupling not blow up below the GUT scale; the numerical difference between the bounds comes from the difference in the evolution of the running couplings in SUSY and the SM. An e^+e^- collider operating at a centre of mass energy ~ 300 GeV would thus be certain¹⁴ to find a Higgs boson if these arguments are valid.

3.2 *MSSM Model Parameters: A Recapitulation*

For the convenience of the reader and for subsequent developments, we first summarize the parameters of the MSSM. In addition to the SM parameters, the MSSM parameters include,

- $\tan\beta$ and the superpotential parameter μ ,
- soft breaking masses for each of the three gauginos: these are given in terms of a single parameter if we assume the gaugino mass unification condition (6).
- There is an independent soft SUSY breaking scalar mass for each SM; *i.e.* $SU(3)\times SU(2)\times U(1)$, matter multiplet. There are thus six slepton masses and nine squark masses for the three families, even if mixing between the generations is ignored.
- Again without inter-generational mixing, there are nine A -parameters, and a B -parameter for the one bilinear term in the MSSM superpotential f_1 .
- Finally, there is the one additional parameter (chosen to be m_{H_p}) that determines the tree-level Higgs boson sector.

We see that without assuming anything more than $SU(3)\times SU(2)\times U(1)$ invariance, the model contains an unmanageably large number of parameters. Assuming grand unification ameliorates the situation to some extent: there are then only two scalar masses per generation of sfermions in $SU(5)$ and only one gaugino mass parameter, but the parameter space is still too large for the phenomenology to be tractable. Inspired by supergravity model studies, many early phenomenological studies assumed that all squarks (sleptons were either assumed to be degenerate with squarks, or to have masses related to $m_{\tilde{q}}$) were degenerate except for D -term splitting. They also incorporated the GUT assumption for gaugino masses. In this case the masses and couplings of all sparticles were determined in terms of relatively few SUSY parameters which were frequently taken to be,

$$m_{\tilde{g}}, m_{\tilde{q}}, m_{\tilde{\ell}}, \mu, \tan \beta, A_t, m_{H_p}. \quad (7)$$

The parameter A_t mainly affects top squark phenomenology, and so, was frequently irrelevant. Other A -terms, being proportional to the light fermion masses, are negligible.

In view of the fact that additional assumptions are necessary, and further, that assumptions based on supergravity models are incorporated into phenomenological analyses, it seems reasonable to explore the implications of these models more seriously. Toward this end, we describe the underlying framework in the following section.

4 Minimal Supergravity Models

When supersymmetry is promoted to a local symmetry, additional fields have to be introduced. The resulting theory¹⁵ which includes gravitation is known as supergravity (SUGRA). It is not our purpose here to study SUGRA models in any detail. In fact, local supersymmetry will not play any direct role in our later considerations. The purpose of this discussion is merely to maintain continuity of development, and also to provide motivation for an economic and elegant framework that has recently become very popular for phenomenological analysis.ⁿ

It was recognised rather early that it is very difficult to construct globally supersymmetric models where SUSY is spontaneously broken at the weak scale. This led to the development of geometric hierarchy models where SUSY is broken in a “hidden” sector at a scale $\mu_s \gg M_W$. This sector is assumed to interact with ordinary particles and their superpartners (the “observable” sector) only via exchange of superheavy particles X . This then suppresses the couplings of the Goldstone fermion (which resides in the hidden sector) to the observable sector: as a result, the effective mass gap in the observable sector is $\mu \sim \frac{\mu_s^2}{M_X}$ which can easily be $\lesssim 1$ TeV even if μ_s is much larger.

An especially attractive realization of this idea stems from the assumption that the hidden and observable sectors interact only gravitationally, so that the scale M_X is $\sim M_{Plank}$. This led to the development of SUGRA GUT models of particle physics. Because supergravity is not a renormalizable theory, we should look upon the resulting Lagrangian, with heavy degrees of freedom integrated out, as an effective theory valid below some ultra-high scale M_X around M_{GUT} or M_{Plank} , in the same way that chiral dynamics describes interactions of pions below the scale of chiral symmetry breaking. Remarkably,

ⁿAlthough this has been discussed by Bagger¹ it seems necessary to include discussion of this important topic for completeness.

this Lagrangian turns out to be just the same as that of a globally supersymmetric $SU(3) \times SU(2) \times U(1)$ model, together with soft SUSY breaking masses and A - and B -parameters of $O(M_{Weak})$.

The economy of the minimal supergravity GUT framework ^o stems from the fact that because of the assumed symmetries, various soft SUSY breaking parameters become related independent of the details of the hidden sector and the low energy effective Lagrangian can be parametrized in terms of just a few parameters. For instance, since the chiral multiplets are universally coupled to the hidden sector (via gravitational interactions), they all acquire the same soft SUSY breaking scalar mass m_0 . Likewise, there is a universal A -parameter, common to all trilinear interactions. The GUT assumption, of course, implies that the soft SUSY breaking gaugino masses are related as in Eq. (6). It should be emphasized that the universality of the scalar masses does not imply that the physical scalar masses of all sfermions are the same. The point is that the parameters in the Lagrangian obtained by integrating out heavy fields should be regarded as renormalized at the high scale M_X at which these symmetries are unbroken. If we use this Lagrangian to compute processes at the 100 GeV energy scale relevant for phenomenology, large logarithms $O(\ln \frac{M_X}{M_W})$ due to the disparity between the two scales invalidate the perturbation expansion. These logarithms can be straightforwardly summed by evolving the Lagrangian parameters down to the weak scale. This is most conveniently done ¹⁶ using renormalization group equations (RGE).

The renormalization group evolution leads to an interesting pattern of sparticle masses, evaluated at the weak scale. ^p For example, gauge boson-gaugino loops result in increased sfermion masses as we evolve these down from M_X to M_W while superpotential Yukawa couplings (which are negligible for the two lightest generations) have just the opposite effect. Since squarks have strong interactions in addition to the electroweak interactions common to all sfermions, the weak scale squark masses are larger than those of sleptons. Neglecting Yukawa couplings, we have to a good approximation,

$$\begin{aligned} m_{\tilde{q}}^2 &= m_0^2 + m_q^2 + (5 - 6)m_{\frac{1}{2}}^2 + D - terms, \\ m_{\tilde{\ell}}^2 &= m_0^2 + m_{\ell}^2 + (0.15 - 0.5)m_{\frac{1}{2}}^2 + D - terms. \end{aligned} \quad (8)$$

^oHere the term minimal refers to the canonical choice of kinetic energy terms for matter and gauge fields. Since supergravity is a non-renormalizable theory, in principle, these terms can arise from higher dimensional operators.

^pThese running masses evaluated at the sparticle mass, or more crudely, at a scale $\sim M_Z$, are not identical to, but are frequently close to the physical masses which are given by the pole of the renormalized propagator ¹⁷.

In Eq. (8), $m_{\frac{1}{2}}$ is the common gaugino mass at the scale m_X . Notice that squarks and sleptons within the same SU(2) doublets are split by the D -terms, once electroweak symmetry is broken. In contrast, various flavours of left- (and separately, right-) type squarks of the first two generations are essentially degenerate, consistent¹⁸ with flavour changing neutral current (FCNC) constraints in the K -meson sector.⁹ The D -terms, which are typically $\leq \frac{1}{2}M_Z^2$, are generally not important when sfermions are heavy. The difference in the coefficients of the $m_{\frac{1}{2}}$ terms reflects the difference between the strong and electroweak interactions alluded to above. Although we have not shown this explicitly, $\tilde{\ell}_R$ which has only hypercharge interactions tends to be lighter than $\tilde{\ell}_L$ as well as $\tilde{\nu}_L$ unless D -term effects are significant. Since $m_{\tilde{g}} = (2.5 - 3)m_{\frac{1}{2}}$, it is easy to see that squark and slepton masses are related by,

$$m_{\tilde{q}}^2 = m_{\tilde{\ell}}^2 + (0.7 - 0.8)m_{\tilde{g}}^2. \quad (9)$$

Here, $m_{\tilde{q}}^2$ and $m_{\tilde{\ell}}^2$ are the squared masses averaged over the squarks or sleptons of the first (or second) generation. In the second term, the unification of gaugino masses has been assumed. Since experimental data, as we will see, requires squarks to be heavier than 150-200 GeV, it immediately follows that the first two generations of squarks are approximately degenerate.

The Yukawa couplings of the top family are certainly not negligible. For very large values of $\tan\beta \sim \frac{m_t}{m_b}$ bottom Yukawa couplings are also important. As mentioned above, these Yukawa interactions tend to reduce the scalar masses at the weak scale. These corrections can overcome the additional m_t^2 in Eq. (8), so that \tilde{t}_L and \tilde{t}_R tend to be significantly lighter than other squarks (of course, by SU(2) invariance, the soft-breaking mass for \tilde{b}_L is the same as that for \tilde{t}_L). In fact, we can say more: because \tilde{t}_R receives top quark Yukawa corrections from both charged and neutral Higgs loops in contrast to \tilde{t}_L which gets corrections just from the neutral Higgs, its squared mass is reduced by (approximately) twice as much as that of \tilde{t}_L . Moreover, as we have already seen, these same Yukawa interactions lead to $\tilde{t}_L - \tilde{t}_R$ mixing, which further depresses the mass of the lighter of the two t -squarks (sometimes referred to as the stop) which we will denote by \tilde{t}_1 . In fact, care must be exercised in the choice of input parameters: otherwise $m_{\tilde{t}_1}^2$ is driven negative, leading to the spontaneous breakdown of electric charge and colour.

The real beauty and economy of this picture comes from the fact that these same Yukawa radiative corrections drive²⁰ electroweak symmetry breaking. Since the Higgs bosons are part of chiral supermultiplets, they also have a

⁹This is a non-trivial observation since alternative mechanisms to suppress FCNC based on different symmetry considerations have been proposed.¹⁹

common mass m_0 at the scale M_X and undergo similar renormalization as doublet sleptons due to gauge interactions; *i.e.* these positive contributions are not very large. The squared mass m_h^2 of the Higgs boson doublet h which couples to the top family, however, receives large negative contributions (thrice those of the \tilde{t}_L squark since there are three different colours running in the loop) from Yukawa interactions, and so can become negative, leading to the correct pattern of gauge symmetry breaking. Furthermore, because $f_t > f_b$, $\tan\beta > 1$. While this mechanism is indeed very pretty, it should probably not be regarded as an explanation of the observed scale of spontaneous symmetry breakdown since it requires that m_0 , the scalar mass at the very large scale M_X be chosen to be ≤ 1 TeV: in other words, the small dimensionless ratio $\frac{m_0}{M_X}$ remains unexplained.

Let us compare the model parameters with our list (7) for the MSSM. Within SUGRA GUTS, we start with GUT scale parameters, m_0 , $m_{\frac{1}{2}}$, A_0 , B_0 and μ_0 . The weak scale parameter μ (actually, μ^2) is adjusted to give the experimental value of M_Z . It is convenient to eliminate B_0 in favour of $\tan\beta$ so that the model is completely specified by just four parameter set (with a sign ambiguity for μ),

$$m_0, m_{\frac{1}{2}}, \tan\beta, A_0, \text{sgn}(\mu), \quad (10)$$

without the need of additional *ad hoc* assumptions as in the MSSM. Comparing with the MSSM parameter set (7) we see that μ and m_{H_p} are no longer free parameters.

SUGRA models lead to a rather characteristic pattern of sparticle masses²¹ and mixings. We have already seen that the first two generations of squarks are approximately degenerate, while the lighter of the t -squarks, and also \tilde{b}_L can be substantially lighter. Also, from Eq. (9) it follows that sleptons may be significantly lighter than the first two generations of squarks if $m_{\tilde{g}} \simeq m_{\tilde{q}}$, and have comparable masses if squarks are significantly heavier than gluinos. We also see that gluinos can never be much heavier than squarks.^r Furthermore, because the top quark is very massive, the value of $|\mu|$ obtained from the radiative symmetry breaking constraint generally tends to be much larger than the electroweak gaugino masses, so that the lighter (heavier) charginos and neutralinos tend to be gaugino-like (Higgsino-like).

It should be kept in mind that while the minimal SUGRA framework provides a very attractive and economic picture, it hinges upon untested as-

^rEllwanger²² has shown that this is a very general result not special to supergravity. It follows from the requirement that $m_{\tilde{q}}^2$ not be driven to negative values below the unification scale, assuming only that there are no new large Yukawa interactions.

assumptions of symmetries about the physics at very high energies. It could be that the GUT assumption is incorrect though this would then require the unification implied²³ by the observed values of gauge couplings at LEP to be purely fortuitous. It could be that the assumption of universal scalar masses (or A -parameters) is wrong. Recall that although we used supergravity couplings between the hidden and observable sectors to argue for this, the common scalar mass was a consequence of the assumed universality of the (gravitational) couplings between the hidden and observable sectors. In other words, the universality of scalar masses is really a result of an assumed global $U(N)$ symmetry of the Lagrangian for transformations amongst the N chiral supermultiplets, an assumption which is, perhaps, reasonable as long as we are near the Planck scale where gravitation presumably dominates GUT gauge interactions. Non-universal masses could result if this $U(N)$ is broken as Bagger has illustrated in his lectures by the explicit introduction of non-renormalizable terms in the superpotential. We should also remember that in the absence of a theory about physics at the high scale, we do not have a really good principle for choosing the scale M_X at which the scalar masses are universal. In practice, most phenomenological calculations set this to be the scale of GUT symmetry breaking where the gauge couplings unify. If, instead, this scale were closer to M_{Planck} the evolution between these scales²⁴ could result in non-universal scalar masses at M_{GUT} : this could have significant impact, particularly on the condition of electroweak symmetry breaking.

Despite these shortcomings, this framework at the very least should be expected to provide a useful guide to our thinking about supersymmetry phenomenology. In spite of the fact that it is theoretically rather constrained, it is consistent with all experimental and even cosmological constraints and even, as we will see, contains a candidate for galactic and cosmological dark matter. But it should be kept in mind that some of the underlying assumptions may prove to be incorrect. For this reason, one should always be careful to test the sensitivity of the phenomenological predictions to the various assumptions, especially when considering the design of future experiments. It is, nevertheless, worth emphasizing that we now have a reasonably flexible yet tractable framework whose underlying assumptions, as we will see, can be subject to direct tests at future colliders.

5 Decays of Supersymmetric Particles

Before we can discuss signatures via which sparticle production might be detectable at colliders, we need to understand how sparticles decay. The conservation of R -parity implies that sparticles can only decay into other sparti-

cles, until the decay cascade terminates in the lightest supersymmetric particle (LSP) which is absolutely stable. There are strong limits²⁵ on the existence of stable or even very long-lived ($\tau > \text{age of the universe}$) coloured or charged sparticles. Such sparticles, which would have been abundantly produced in the Big Bang, would bind to ordinary particles to form exotic atoms or nuclei²⁶. For masses up to 1 TeV, their expected abundances are in the range $\gtrsim 10^{-10}$, whereas we know experimentally that these abundances are $< \mathcal{O}(10^{-12} - 10^{-29})$ depending on the new particle mass and also on the element whose exotic isotope is being searched for.

Within the MSSM, the null result of these searches is taken to imply that the LSP must be a weakly interacting neutral particle; *i.e.* it must be either the lightest neutralino \tilde{Z}_1 , or one of the sneutrinos. We will see later that sneutrinos are excluded as the LSP if we also require that they make up galactic dark matter. Within the SUGRA framework, the LSP could also be the gravitino — the SUSY partner of the graviton. Unless it is extremely light, it couples to other particles with gravitational strength couplings, so that it is effectively decoupled for the purposes of collider phenomenology: then, the next lightest SUSY particle, which will decay outside the detector, plays the role of the LSP.^s In this case (or if R -parity is not conserved), however, the “effective” (or actual) LSP may even be charged or coloured. Throughout most of these lectures, we will assume that \tilde{Z}_1 is the LSP.

We note here that regardless of details the neutral LSP’s which are produced at the termination of the SUSY decay cascade behave like stable, heavy neutrinos in the experimental apparatus in that they escape without depositing any energy. Thus apparent missing energy (\cancel{E}) and an imbalance of transverse momentum (\cancel{p}_T) are generally regarded as canonical signatures of supersymmetry.

5.1 Sfermion Decays

We have seen in Sec. 2 that gauginos and Higgsinos couple sfermions to fermions. Since we have also assumed that \tilde{Z}_1 is the LSP, the decay $\tilde{f}_{L,R} \rightarrow f \tilde{Z}_1$ ($f \neq t$) is always allowed. Depending on sparticle masses, the decays

$$\tilde{f}_{L,R} \rightarrow f \tilde{Z}_i, \tilde{f}_L \rightarrow f' \widetilde{W}_i \quad (11)$$

to other neutralinos or to charginos may also be allowed. The chargino decay modes of \tilde{f}_R only proceed via Yukawa interactions, and so are negligible for all

^sAs long as the next lightest sparticle is neutral, SUSY phenomenology at colliders is essentially unaltered. The late decay of this effective LSP can potentially spoil the successful predictions of Big Bang nucleosynthesis as discussed by Moroi.²⁷

but t -squarks. Unlike sleptons, squarks also have strong interactions, and so can also decay into gluinos via,

$$\tilde{q}_{L,R} \rightarrow q\tilde{g}, \quad (12)$$

if $m_{\tilde{q}} > m_{\tilde{g}}$. Unless suppressed by phase space, the gluino decay mode of squarks dominates, so that squark signatures are then determined by the decay pattern of gluinos. If $m_{\tilde{q}} < m_{\tilde{g}}$, squarks, like sleptons, decay^{28,29} to charginos and neutralinos. The important thing to remember is that sfermions dominantly decay via the two-body mode.

The various partial decay widths can be easily computed using the Lagrangian we have described above. Numerical results may be found in the literature for both sleptons³⁰ and squarks²⁹ and will not be repeated here. The following features, however, are worthy of note:

- The electroweak decay rates are $\sim \alpha m_{\tilde{f}}$ corresponding to lifetimes of about $10^{-22}(\frac{100 \text{ GeV}}{m_{\tilde{f}}})$ seconds. Thus sfermions decay without leaving any tracks in the detector. We will leave it to the reader to check that the same is true for the decays of other sparticles discussed below.
- Light sfermions directly decay to the LSP. For heavier sfermions, other decays also become accessible. Decays which proceed via the larger SU(2) gauge coupling are more rapid than those which proceed via the smaller U(1) coupling (Higgs couplings are negligible). Thus, for \tilde{f}_L , the decays to charginos dominate unless they are kinematically suppressed, whereas \tilde{f}_R ($f \neq t$) mainly decays into the neutralino with the largest U(1) gaugino component.
- Very heavy sleptons (and squarks, if the gluino mode is forbidden) preferentially decay into the lighter (heavier) chargino (\tilde{f}_L only) and the lighter neutralinos $\tilde{Z}_{1,2}$ (the heavier neutralinos $\tilde{Z}_{3,4}$) if $|\mu|$ ($m_{\tilde{g}}$) is very large. This is because $\tilde{W}_1, \tilde{Z}_{1,2}$ ($\tilde{W}_2, \tilde{Z}_{3,4}$) are the sparticles with the largest gaugino components.

Top Squark Decays: We have seen that t -squarks are special in that (i) the mass eigenstates are parameter-dependent mixtures of \tilde{t}_L and \tilde{t}_R , (ii) \tilde{t}_1 , the lighter of the two states may indeed be much lighter than all other sparticles (except, of course, for phenomenological reasons, the LSP) even when other squarks and gluinos are relatively heavy, and (iii) top squarks couple to charginos and neutralinos also via their Yukawa components. As a result the decay patterns of \tilde{t}_1 can differ considerably from those of other squarks.

The decay $\tilde{t}_1 \rightarrow t\tilde{g}$ will dominate as usual if it is kinematically allowed. Otherwise, the decays to charginos and neutralinos, if allowed, form the main decay modes. Since m_t is rather large, it is quite possible that the decay $\tilde{t}_1 \rightarrow t\tilde{Z}_1$ is kinematically forbidden, and $\tilde{t}_1 \rightarrow b\tilde{W}_1$ is the only tree-level two body decay mode that is accessible, in which case it will obviously dominate. If the stop is lighter than $m_{\tilde{W}_1} + m_b$, and has a mass smaller than about 125 GeV (which, we will see, is in the range of interest for experiments at the Tevatron), the dominant decay mode of \tilde{t}_1 comes from the flavour-changing $\tilde{t}_1 - \tilde{c}_L$ loop level mixing induced by weak interactions³¹ and the decay $\tilde{t}_1 \rightarrow c\tilde{Z}_1$ dominates the allowed (at least four-body) tree level decays. If $m_{\tilde{t}_1} \sim 175 - 225$ GeV, the three-body decays $\tilde{t}_1 \rightarrow bW\tilde{Z}_1$ may be accessible, with the two body decays $\tilde{t}_1 \rightarrow b\tilde{W}_1$ and $\tilde{t}_1 \rightarrow t\tilde{Z}_1$ still closed. How this decay, which could be of interest for stop searches at Tevatron upgrades or at future e^+e^- linear colliders or the Large Hadron Collider, compares with the loop decay is currently under investigation.

5.2 Gluino Decays

Since gluinos have only strong interactions, they can only decay via

$$\tilde{g} \rightarrow \bar{q}\tilde{q}_{L,R}, q\tilde{q}_{L,R}, \quad (13)$$

where the squark may be real or virtual depending on squark and gluino masses. If $m_{\tilde{g}} > m_{\tilde{q}}$, \tilde{q}_L and \tilde{q}_R are produced in equal numbers in gluino decays (except for phase space corrections from the non-degeneracy of squark masses). In this case, since \tilde{q}_R only decays to neutralinos, neutralino decays of the gluino dominate. If, as is more likely, $m_{\tilde{g}} < m_{\tilde{q}}$, the squark in Eq. (13) is virtual and decays via Eq. (11), so that gluinos decay via three body modes,

$$\tilde{g} \rightarrow q\bar{q}\tilde{Z}_i, q\bar{q}'\tilde{W}_i. \quad (14)$$

In contrast to the $m_{\tilde{g}} > m_{\tilde{q}}$ case, gluinos now predominantly decay^{28,29} into charginos because of the large SU(2) gauge coupling, and also into the neutralino with the largest SU(2) gauge component. For small values of μ ($\ll \mu_2$), these may well be the heavier chargino and the heaviest neutralino²⁹; if instead μ is relatively large, as is generally the case in SUGRA type models, the \tilde{W}_1 and \tilde{Z}_2 decays of gluinos frequently dominate.

We should also point out that our simplistic discussion above neglects differences between various squark masses. As we have seen in the last section, however, third generation squarks \tilde{t}_1 and $\tilde{b}_1 \sim \tilde{b}_L$ may in fact be substantially lighter than the other squarks. It could even be³² that $\tilde{g} \rightarrow \bar{b}\tilde{b}_1$ and/or

$\tilde{g} \rightarrow t\bar{t}_1$ are the only allowed two-body decays of the gluino in which case gluino production will lead to final states with very large b -multiplicity, and possibly also hard, isolated leptons from the decays of top or stop quarks. Even if these decays are kinematically forbidden, decays to third generation fermions may nonetheless be large because of enhancement of the \tilde{t}_1 and \tilde{b}_1 propagators (recall that the decay rates roughly depend on $\frac{1}{m_q^4}$) with qualitatively the same effect.

Finally, we note that there are some regions of parameter space where the radiative decay,

$$\tilde{g} \rightarrow g\tilde{Z}_i, \quad (15)$$

can be important³³. This decay, which occurs via third generation squark and quark loops, is typically enhanced relative to the tree-level decays if the neutralino contains a large \tilde{h} component (which has large Yukawa couplings to the top family).

5.3 Chargino and Neutralino Decays

Within the MSSM framework where baryon and lepton number are conserved, charginos and neutralinos can either decay into lighter charginos and neutralinos and gauge or Higgs bosons, or into fermion-sfermion pairs if these decays are kinematically allowed. We will leave it as an exercise to the reader to make a listing of all the allowed modes and refer the reader to the literature^{10,30} for various formulae and numerical values of the branching fractions. If these two-body decay modes are all forbidden, the charginos and neutralinos decay via three body modes,

$$\widetilde{W}_i \rightarrow f\bar{f}'\tilde{Z}_j, \widetilde{W}_2 \rightarrow f\bar{f}\widetilde{W}_1 \quad (16)$$

$$\tilde{Z}_i \rightarrow f\bar{f}\tilde{Z}_j \text{ or } f\bar{f}'\widetilde{W}_1, \quad (17)$$

mediated by virtual gauge bosons or sfermions (amplitudes for Higgs boson mediated decays, being proportional to fermion masses are usually negligible). Typically, only the lighter chargino and the neutralino \tilde{Z}_2 decay via three body modes, since the decays $\tilde{Z}_{3,4} \rightarrow \tilde{Z}_1 Z$ or $\tilde{Z}_1 H_\ell$ and $\widetilde{W}_2 \rightarrow W\tilde{Z}_1$ are often kinematically accessible. Of course if the \tilde{Z}_2 or \widetilde{W}_1 are heavy enough they will also decay via two body decays: these decays of \tilde{Z}_2 are referred to as “spoiler modes” since, as we will see, they literally spoil^t the clean leptonic signal via which \tilde{Z}_2 may be searched for.

^tThe decay to Higgs does not yield leptons, whereas the decay to Z has additional backgrounds from SM Z sources.

For sfermion masses exceeding about M_W , W -mediated decays generally dominate the three body decays of \widetilde{W}_1 , so that the leptonic branching for its decays fraction is essentially the same as that of the W ; *i.e.* 11% per lepton family. An exception occurs when μ is extremely large so that the LSP is mainly a U(1) gaugino and \widetilde{W}_1 dominantly an SU(2) gaugino. In this case, the $W\widetilde{W}_1\widetilde{Z}_1$ coupling is considerably suppressed: then, the amplitudes for \widetilde{W}_1 decays mediated by virtual sfermions may no longer be negligible, even if sfermions are relatively heavy, and the leptonic branching fractions may deviate substantially from their canonical value of 11%.

One may analogously expect that \widetilde{Z}_2 decays are dominated by (virtual) Z^0 exchange if sfermion masses substantially exceed M_Z . This is, however, not true since the Z^0 couples only to the Higgsino components of the neutralinos, so that if either of the neutralinos in the decay $\widetilde{Z}_2 \rightarrow \widetilde{Z}_1 f \bar{f}$ has small Higgsino components the Z^0 contribution may be strongly suppressed, and the contribution from amplitudes involving relatively heavy sfermions may be non-negligible. This phenomenon is common in SUGRA models where $|\mu|$ is generally much larger than the electroweak gaugino masses, and \widetilde{Z}_1 and \widetilde{Z}_2 are, respectively, mainly the hypercharge and SU(2) gauginos. If, in addition, $m_{\tilde{q}} \sim m_{\tilde{g}}$, we see from Eq. (9) that sleptons are much lighter than squarks, so that the leptonic decays $\widetilde{Z}_2 \rightarrow \ell \bar{\ell} \widetilde{Z}_1$, which lead to clean signals at hadron colliders, may be considerably enhanced.³⁴ There are, however, other regions of parameter space, where sleptons are relatively light, but the amplitudes from virtual slepton exchanges interfere destructively with the Z^0 mediated amplitudes, and lead to a strong suppression of this decay.^{35,36} Of course, the branching fraction for the three-body decay is tiny if two-body “spoiler modes” $\widetilde{Z}_2 \rightarrow Z \widetilde{Z}_1$ or $\widetilde{Z}_2 \rightarrow H_t \widetilde{Z}_1$ are kinematically allowed. For basically the same reasons the decay $\widetilde{Z}_2 \rightarrow \widetilde{W}_1 f \bar{f}'$ which is mediated by virtual W exchange, even though it is kinematically disfavoured, can sometimes be competitive³⁷ with the LSP decay mode of \widetilde{Z}_2 . A complete set of formulae useful for evaluating the rates for the three body decays of charginos and neutralinos may be found^u in Bartl *et. al.*³⁸

Finally, we note that there are regions of parameter space where the rate for the two body radiative decay

$$\widetilde{Z}_2 \rightarrow \widetilde{Z}_1 \gamma \quad (18)$$

which is mediated by $f \bar{f}$ and gauge boson-gaugino loops may be comparable^{39,40} to that for the three body decays. These are important in two different

^uWe warn the reader that their notation and conventions do not match those used in these lectures, so some care should be exercised in transcribing these into a common notation.

cases: (i) if one of the neutralinos is photino-like and the other Higgsino-like, both Z^0 and sfermion mediated amplitudes are small since the photino (Higgsino) does not couple to the Z -boson (sfermion), and (ii) both neutralinos are Higgsino-like and very close in mass (this occurs for small values of $|\mu|$); the strong suppression of the three-body phase space then favours the two-body decay. We mention here that neither of these cases is particularly likely, especially within the SUGRA framework.

5.4 Higgs Boson Decays

Unlike in the SM, there is no clear dividing line between the phenomenology of sparticles and that of Higgs bosons, since as we have just seen, Higgs bosons can also be produced via cascade decays of heavy sparticles. Higgs boson decay patterns exhibit¹⁰ a complex dependence on model parameters. Unfortunately, we will not have time to discuss these here, and we can only refer the reader to the literature. We will, therefore, confine ourselves to mentioning a few points that will be important for later discussion.

In SUGRA models, all but the lightest Higgs scalar tend to be (but are not always) rather heavy and so are not significantly produced either in sparticle decay cascade decays or directly at colliders. Within the more general MSSM framework, the scale of their masses is fixed by m_{H_p} , which is an independent parameter. In this limit, H_ℓ which has a mass smaller than ~ 130 GeV, behaves like the SM Higgs boson, while H_h , H_p and H^\pm are approximately decoupled from vector boson pairs. The phenomenology is then relatively simple: the decay $H_\ell \rightarrow b\bar{b}$ which occurs via b -quark Yukawa interactions dominates, unless charginos and/or neutralinos are also light; then, decays of H_ℓ into neutralino or chargino pairs, which occur via the much larger gauge coupling, may be dominant. The invisible decay $H_\ell \rightarrow \tilde{Z}_1\tilde{Z}_1$, is clearly the one most likely to be accessible, and has obvious implications for Higgs phenomenology. These supersymmetric decay modes are even more likely for the heavier Higgs bosons, particularly if their decay to $t\bar{t}$ pairs is kinematically forbidden; this is especially true for H_p which cannot decay to vector boson pairs, but also for H_h since its coupling to VV pairs ($V = W, Z$) is suppressed for $m_{H_h} \gtrsim 200$ GeV. The decays $H_p \rightarrow H_\ell Z$ and $H_h \rightarrow H_\ell H_\ell$ can, of course, be important, while $H_h \rightarrow H_p H_p$ is usually inaccessible. Finally, charged Higgs bosons H^\pm mainly decay via the $t\bar{b}$ mode unless this channel is kinematically forbidden. Then, they mainly decay via $H^\pm \rightarrow WH_\ell$, or if this is also kinematically forbidden, via $H^\pm \rightarrow c\bar{s}$ or $H^\pm \rightarrow \bar{\tau}\nu$ with branching fractions depending on $\tan\beta$.

6 Sparticle Production at Colliders

Since R -parity is assumed to be conserved, sparticles can only be pair produced by collisions of ordinary particles. At e^+e^- colliders sparticles (such as charged sleptons and sneutrinos, squarks and charginos) with significant couplings to either the photon or the Z -boson can be produced via s -channel γ and Z processes, with cross sections comparable with $\sigma(e^+e^- \rightarrow \mu^+\mu^-)$, except for kinematic and statistical factors. Selectron and electron sneutrino production may also occur via t -channel neutralino and chargino exchange, while sneutrino exchange in the t -channel will contribute to chargino pair production. Neutralino production, which proceeds via Z exchange in the s -channel and selectron exchange in the t and u channels, may be strongly suppressed if the neutralinos are gaugino-like and selectrons are relatively heavy. Cross section formulae as well as magnitudes of the various cross sections may be found *e.g.* in Baer *et. al.*³⁰

At hadron colliders, the situation is somewhat different. Since sparticle production is a high Q^2 process, the underlying elementary SUSY process is the inelastic collision of quarks and gluons inside the proton.⁴¹ In other words, it is the partonic cross section that is computable within the SUSY framework. This cross section is then convoluted with parton distribution functions to obtain the inclusive cross section for SUSY particle production. Thus, unlike at electron-positron colliders, only a fraction of the total centre of mass energy is used for sparticle production. The balance of the energy is contained in the underlying low p_T event which only contaminates the high p_T signal of interest.

Squarks and gluinos, the only strongly interacting sparticles, have the largest production cross sections unless their production is kinematically suppressed. These cross sections⁴² are completely determined in terms of their masses by QCD and do not depend on the details of the supersymmetric model. QCD corrections to these have also been computed.⁴³ Squarks or gluinos can be also be produced⁴⁴ in association with charginos or neutralinos via diagrams involving one strong and one electroweak vertex. Finally, \tilde{W}_i and \tilde{Z}_j can be produced by $q\bar{q}$ annihilation via processes with W or Z exchange in the s -channel, or squark exchange in the t (and, for neutralino pairs only, also the u) channel.

The cross sections for various processes at a 2 TeV $p\bar{p}$ collider (corresponding to the Main Injector (MI) upgrade of the Tevatron) are illustrated in Fig. 1, while those for a 14 TeV pp collider (the recently approved LHC) are shown in Fig. 2. We have illustrated our results for (a) $m_{\tilde{q}} = m_{\tilde{g}}$, and (b) $m_{\tilde{q}} = 2m_{\tilde{g}}$ and fixed other parameters at some representative values shown. These figures

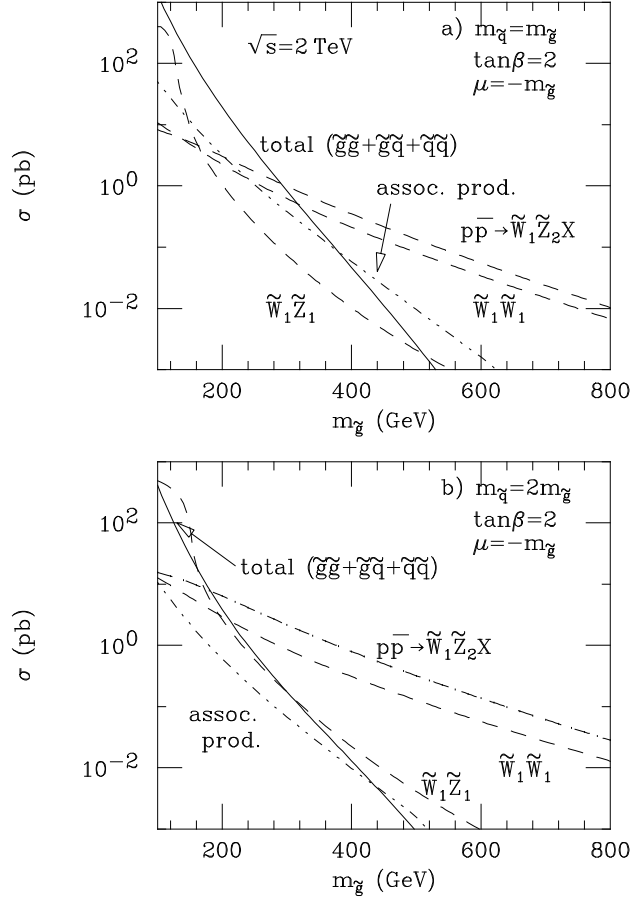


Figure 1: Total cross sections for various sparticle production processes by $p\bar{p}$ collisions at $\sqrt{s} = 2$ TeV.

help us decide what to search for. While squarks and gluinos are the obvious thing to focus the initial search on, we see from Fig. 1 that at even the MI (and certainly at the TeV33 upgrade being envisioned for the future), the maximal reach is likely to be obtained via the electroweak production of charginos and neutralinos, provided of course that their decays lead to detectable signals.^v In contrast, we see from Fig. 2 that gluino (and, possibly, squark) production processes offer the best opportunity for SUSY searches at the LHC for gluino masses up to 1 TeV (recall that this is roughly the bound from fine-tuning considerations⁴⁵) even if squarks are very heavy.

7 Simulation of Supersymmetry Events

Once produced, sparticles rapidly decay into other sparticles until the decay cascade terminates in a stable LSP. It is the end products of these decays that are detectable in experiments and which, it is hoped, can provide experimental signatures for supersymmetry. The evaluation of these signatures obviously entails a computation of the branching fractions for the decays of all the sparticles, and further, keeping track of numerous cascade decay chains for every pair of parent sparticles. Many groups have generated computer programs to calculate these decay processes. For any set of MSSM parameters (7) (or alternatively, for a SUGRA parameter set (10)), a public access program known as ISASUSY (ISASUGRA) which can be extracted from the Monte Carlo program ISAJET⁴⁶ lists all sparticle and Higgs boson masses as well as their decay modes along with the corresponding partial widths and branching fractions.

Event generator programs provide the link between the theoretical framework of SUSY which provides, say, cross sections for final states with quarks and leptons, and the long-lived particles such as π , K , γ , e , μ *etc.* that are ultimately detected in real experiments. Many groups have combined sparticle production and decay programs to create parton level event generators which may be suitable for many purposes. More sophisticated generators include other effects such as parton showers, heavy flavour decays, hadronization of gluons and quarks, a model of the underlying event, *etc.* These improvements have significant impact upon detailed simulations of, for instance, the jets plus isolated multi-lepton signal from squark and gluino production at the LHC.

These lectures are not the place to discuss these generators in detail, so we will content ourselves with providing some information of what is available today. ISAJET 7.14 is probably the most comprehensive, but by no means complete, SUSY generator for simulation at hadron colliders. Mrenna⁴⁷ has,

^vThis conclusion crucially depends on the validity of the gaugino mass unification condition Eq. (6).

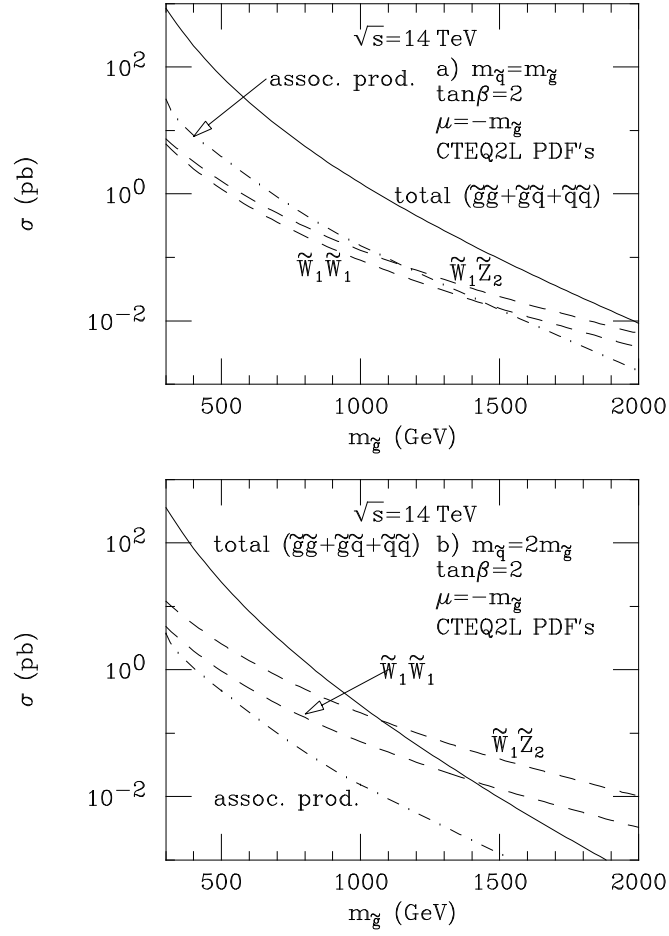


Figure 2: Total cross sections for various particle production processes by pp collisions at $\sqrt{s} = 14$ TeV.

very recently, compiled an independent event generator for SUSY simulation at hadron colliders. ISAJET and SUSYGEN⁴⁸ are two general purpose generators available for simulation of supersymmetry at e^+e^- colliders that include all $2 \rightarrow 2$ production processes and cascade decays of all sparticles. SUSYGEN includes initial state photon radiation, and can be interfaced to the LUND JETSET string hadronization program. Neither of these generators currently incorporates spin correlations or polarization of the incoming beams (due to be included in the next release of ISAJET), while final state decay matrix elements are included only in SUSYGEN. Specialized generators^{49,50,51} that remedy these defects are also available, but these can only be used for the simulation of specific SUSY reactions with specific final states.

8 Observational Constraints on Supersymmetry

The non-observation of any supersymmetric signal at either LEP⁵² or at the Tevatron^{53,54} provide the most direct lower limits on sparticle masses. Indirect limits may also come from virtual effects of SUSY particles on rare processes (*e.g.* flavour changing neutral currents or proton decay) or from cosmological considerations such as an over-abundance of LSP's resulting in a universe that would be younger than the age of stars. While these indirect limits can be important, they are generally sensitive to the details of the model: the non-observation of loop effects could be a result of accidental cancellation with some other new physics loops (so care must be exercised in extracting limits on sparticle masses), proton decay⁵⁵ is somewhat sensitive to assumptions about GUT scale physics while the cosmological constraints⁵⁶ can be simply evaded by allowing a tiny violation of R -parity conservation which would have no impact on collider searches. We should stress that we do not mean to belittle these constraints which lead to important bounds in any *given* framework (for instance, minimal SUGRA SU(5)), but should also recognise that these bounds are likely to be more model-dependent than direct constraints from collider experiments. It is, however, only for reasons of time that we will mainly confine ourselves to direct limits from colliders.

The cleanest limits on sparticle masses come from experiments at LEP. The agreement⁵⁷ of Γ_Z with its expectation in the SM gives⁵⁸ essentially model-independent lower limits of 30-45 GeV on the masses of charginos, squarks, sneutrinos and charged sleptons whose couplings to Z^0 were fixed by gauge symmetry. These limits^w do not depend on how sparticles decay. Likewise,

^wThe same considerations also exclude spontaneous R -violation via a vev of a doublet sneutrino because the associated Goldstone boson sector would then have gauge couplings to Z^0 and make too large a contribution to Γ_Z .

the measurement of the invisible width of the Z^0 which gives the well-known bound on the number of light neutrino species, yields a lower limit on $m_{\tilde{\nu}}$ only 2-5 GeV below $\frac{M_Z}{2}$ if the sneutrino decays invisibly via $\tilde{\nu} \rightarrow \nu \tilde{Z}_1$, even if only one of the sneutrinos is light enough to be accessible in Z^0 decays.^x In contrast, the bounds on neutralino masses are very sensitive to the model parameters because for large $|\mu|$, as we have already pointed out, the neutralino may be dominantly a gaugino with strongly suppressed couplings to the Z^0 .

LEP experimentalists also perform direct searches for sparticles whose decays frequently lead to extremely characteristic final states.⁶¹ For instance, slepton (squark) pair production followed by the direct decay of the sfermion to the LSP leads to a pair of hard, acollinear leptons (jets) together with \cancel{p}_T . Chargino production can lead to events with acollinear jet pairs, a lepton + jet + \cancel{p}_T and also acollinear leptons + \cancel{p}_T . Such event topologies are very distinctive and do not occur in the SM. Thus the observation of just a handful of such events would suffice to signal new physics. The non-observation of SUSY signals in LEP experiments thus implies a lower bound on the masses of sfermions and charginos very close to the kinematic limit. The reactions $e^+e^- \rightarrow \tilde{Z}_1\tilde{Z}_2$, $\tilde{Z}_2\tilde{Z}_2$ can also lead to similarly characteristic final states. As explained above, non-observation of such signals do not lead to bounds on neutralino masses, but do serve to exclude regions of SUSY parameter space.⁶²

Although LEP experiments have resulted in a lower bound $\sim \frac{M_Z}{2}$ on the squark mass, the search for strongly interacting sparticles is best carried out at hadron colliders by searching for \cancel{E}_T events from $\tilde{q}\tilde{q}$, $\tilde{g}\tilde{q}$ and $\tilde{g}\tilde{g}$ production. The final states from the cascade decays of gluinos and squarks leads to events consisting of several jets plus possibly leptons and \cancel{E}_T . For an integrated luminosity of about 10-20 pb^{-1} on which the analyses of the Run IA of the CDF and D0 experiments are based, the classic \cancel{E}_T channel offers the best hope for detection of supersymmetry. The non-observation of \cancel{E}_T events above SM background expectations (after cuts to increase the signal relative to background) has enabled the D0 collaboration⁵⁴ to infer a limit of 173 GeV on $m_{\tilde{q}}$, improving on the published CDF limit of about 100 GeV. The region of the $m_{\tilde{q}} - m_{\tilde{g}}$ plane excluded by these analyses depends weakly on other SUSY parameters and is shown in Fig. 3 for $\mu = -250$ GeV and $\tan\beta = 2$. We see that

^xExperiments searching for neutrino-less double beta decay can detect the recoil of the nucleus. If stable sneutrinos are the LSP and their density is large enough to form all of the galactic dark matter their flux would be high enough to be detectable via elastic scattering from nuclei in these experiments. As a result, sneutrinos with masses between 12-20 GeV and about 1 TeV are excluded.⁵⁹ The Kamiokande experiment⁶⁰, from a non-observation of high energy solar neutrinos produced by the annihilation of gravitationally accumulated sneutrinos in the sun exclude $3 \text{ GeV} \leq m_{\tilde{\nu}} \leq 25 \text{ GeV}$. These limits, when combined with the LEP bounds clearly disfavour the sneutrino as the stable LSP.

Squark & Gluino Exclusion Contour

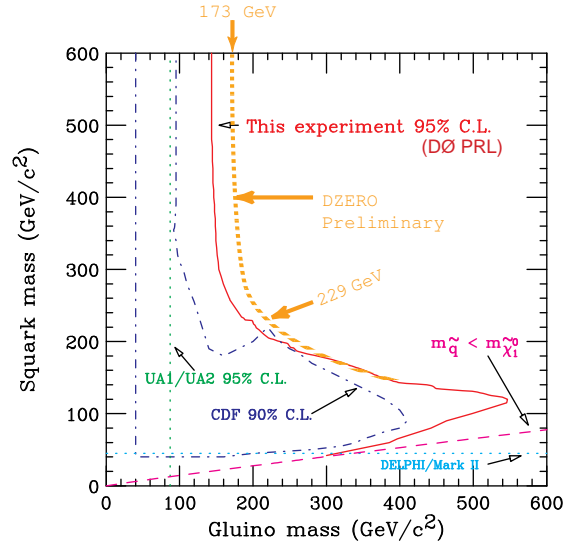


Figure 3: Regions of the $m_{\tilde{g}}$ vs. $m_{\tilde{q}}$ plane excluded by searches for $E_T + jets$ events at various colliders, for $\tan\beta = 2$ and $\mu = -250$ GeV. This figure shows the latest results from the D0 experiment and is taken from Claes.⁵⁴

the lower bound on the mass improves to 229 GeV if squarks and gluinos have the same mass. Since then, the CDF and D0 experiments have collectively accumulated about 150 pb^{-1} of integrated luminosity, and should begin to be sensitive to interesting multilepton signatures which we will discuss when we address prospects for SUSY searches in the future.

Before closing this section, we briefly remark about potential constraints from “low energy” experiments, keeping in mind that these may be sensitive to model assumptions. As discussed in the lectures by Hewett⁶³ the measurements of the inclusive $b \rightarrow s\gamma$ decay by the CLEO experiment⁶⁴ and its agreement with SM expectations constrain the *sum* of SUSY contributions to this process.^y Supersymmetry also allows for new sources of CP violation⁶⁶ in gaugino masses or A -parameters. These phases, which must be smaller than $\sim 10^{-3}$ in order that the electric dipole moment of the neutron not exceed its experimental bound, are set to zero in the MSSM. We will leave it to Hewett to discuss the implications of these novel sources of CP violation.

Finally, we note that because SUSY, unlike technicolour, is a decoupling theory,²³ the agreement of the LEP data with SM expectations is not hard to accommodate. We just have to make the sparticles heavier than 100-200 GeV. But by the same token, it is not easy⁶⁷ to accommodate the observed deviation in the value of $R_b = \Gamma(b\bar{b})/\Gamma(Z \rightarrow \text{hadrons})$ (and even more so⁶⁸ for R_c). While the data appear to prefer a light \tilde{t}_1 and a light chargino, or a light H_p with large $\tan\beta$, it seems hard to obtain a large enough effect to explain the “anomalies”.

9 Searching for Supersymmetry at Future Colliders and Supercolliders

9.1 e^+e^- Colliders

LEP is scheduled to enter its second phase around the end of 1995. The energy of LEP2 is initially expected to be about 140 GeV, but soon should be increased to beyond the WW threshold. The signals for sparticles are much the same as discussed in the last section. The significant difference is that while SM backgrounds can be easily removed below the WW threshold, the separation of the SUSY signal from W -pair production requires more effort. This should not be very surprising since the W is a heavy particle and its decays can lead to both acollinear dilepton + \cancel{E}_T as well as $jets + \ell + \cancel{E}_T$ and $jets + \cancel{E}_T$ event topologies. Another possible complication to be kept in mind as we search for heavier sparticles is that cascade decay channels may begin to open up. This

^yA very recent analysis⁶⁵ suggests that the theoretical error is about twice that assumed in many analyses; this will somewhat relax the restrictions that have been claimed in the literature.

should not pose too much of a problem, however, since the energy is expected to be increased in stages. Thus, for example, one may expect to see chargino production before the production of sleptons which are heavy enough to decay to charginos sets in.

Signals for sparticle production at LEP2 have been studied in great detail^{61,69,70} assuming that sparticles decay directly to the LSP. Below the WW threshold, they are readily detectable in exactly the same way as at LEP. Above that, the production of W^+W^- pairs, which has a very large cross section ~ 18 pb (compared to 0.2 pb for smuons and ~ 10 pb for charginos with mass about M_W) is a formidable background. The situation is not as bad as it may appear on first sight. For WW events to fake sleptons, both W 's have to decay to the particular flavour of leptons, which reduces background by two orders of magnitude. Further rejection of background may be obtained by noting that while slepton events are isotropic, the leptons from W decay exhibit strong backward-forward asymmetry. Thus by selecting from the sample of acollinear $\mu^+\mu^-$ events those events where the fast muon in the hemisphere in the e^- beam direction has the opposite sign to that expected from a muon from W decay, it is possible to reduce the background by a factor of five, with just 50% loss of signal.

The strategy for charginos is more complicated⁷⁰ and will not be detailed here. We will only mention that here the clean environment of electron-positron colliders plays a crucial role. The idea is to make use of the kinematic differences between the two-body decay of the W into a massless neutrino, and the three body decay of the chargino into the massive LSP. Using the cuts detailed in Ref.⁷⁰, it should be possible to detect charginos up to within a few GeV from the kinematic limit in the mixed lepton plus jet channel. Neutralino signals, as we should by now anticipate, are sensitive to model parameters. A recent analysis⁷¹ within the framework of the SUGRA models describes strategies to optimize these signals, and also separate them from other SUSY processes.

Higher energy electron-positron colliders will almost certainly be linear colliders, since synchrotron radiation loss in a circular machines precludes the possibility of increasing the machine energy significantly beyond that of LEP2. Several laboratories are evaluating the prospects for construction of a 300-500 GeV collider, whose energy may later be increased to 1 TeV, or more: these include the Next Linear Collider (NLC) program in the USA, the Japanese Linear Collider (JLC) program in Japan, the TESLA and CLIC programs in Europe, and VLEPP in the former Soviet Union. The search for the lightest charged sparticle, be it the chargino or the slepton (or perhaps the \tilde{t}_1) should proceed along the same lines^{72,51,73} as at LEP2 and discovery should be possible essentially all the way to the kinematic limit. Of course, because production

cross sections rapidly decrease with energy, a luminosity of 10-30 fb^{-1}/yr will be necessary. For the more massive sparticles, cascade decays need to be incorporated, and only relatively preliminary work³⁰ exists on this. Nonetheless, all indications are that at these facilities discovery of any massive new particles with electroweak couplings will not pose serious difficulties. A machine with a centre of mass energy of about 700-1000 GeV should be able to search for charginos up to 350-500 GeV, and so, assuming the gaugino mass unification condition, will cover the parameter space of weak scale supersymmetry.

It is also worth mentioning that one can exploit the availability of polarized beams to greatly reduce SM backgrounds: for example, the cross section for WW production which is frequently the major background is tiny for a right-handed electron beam. While the availability of polarized beams and the clean environment of e^+e^- collisions clearly facilitates the extraction of the signal, we will see later that these capabilities play a really crucial role for the determination of sparticle properties which, in turn, serves to discriminate between models.

Before closing, we should mention that e^+e^- colliders are ideal facilities to search for Higgs bosons.⁷⁴ At LEP2, one can typically search for Higgs bosons with a mass up to about $\sqrt{s} - 100$ GeV; An e^+e^- collider operating at 300 GeV would be virtually guaranteed to find one of the Higgs bosons if the MSSM framework, with its weakly coupled Higgs sector, is correct, although it may not be possible to distinguish this from the Higgs boson of the SM. In contrast, we will see that at the LHC, the discovery of the Higgs boson cannot be guaranteed even with relatively optimistic (but not unrealistic) assumptions about detector capabilities.

9.2 Future Searches at Hadron Colliders

Tevatron Upgrades

The CDF and D0 experiments have together already collected an integrated luminosity of about 150 pb^{-1} and are each expected to accumulate $\sim 100 pb^{-1}$ by the end of the current run, to be compared with 10-20 pb^{-1} for the data set on which the \cancel{E}_T analyses described in the last section were based. It is thus reasonable to explore whether an analysis of this data can lead to other signatures for supersymmetry. Of course, the size of the data sample will increase by yet another order of magnitude after about a year of MI operations and, by significantly more, if the TeV33 upgrade, with its design luminosity of $\sim 10 fb^{-1}/yr$, comes to pass.

Gluinos and Squarks: While the increase in the data sample will obviously result in an increased reach via the \cancel{E}_T channel⁷⁵, we have already seen that

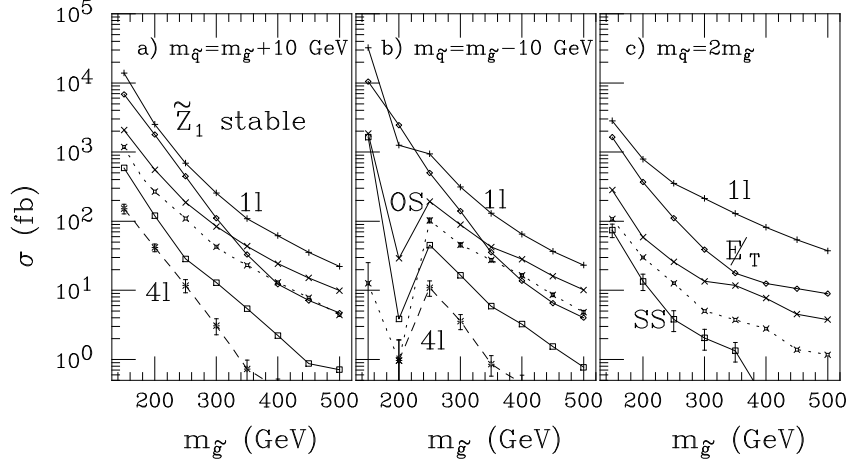


Figure 4: Cross sections (in fb) at the Tevatron ($\sqrt{s} = 1.8$ TeV) for various event topologies after cuts described in Ref. ⁷⁹ from which this figure is taken. We take $\mu = -m_{\tilde{g}}$, $\tan \beta = 2$, $A_t = A_b = -m_{\tilde{g}}$ and $m_{H_p} = 500$ GeV. The \cancel{E}_T cross sections are labelled with diamonds, the $1\text{-}\ell$ cross sections with crosses, the $\ell^+\ell^-$ cross sections with x's and the SS ones with squares. The dotted curves are for the 3ℓ cross sections while the dashed curves show the cross sections for 4ℓ events. For clarity, error bars are shown only on the lowest lying curve; on the other curves, these error bars are significantly smaller. We note that the $m_{\tilde{g}} = 150$ GeV case in b) is already excluded by the LEP constraints on the Z width, since in this case the sneutrino mass is just 26 GeV.

the cascade decays of gluinos and squarks lead to novel signals (n jets plus m leptons plus \cancel{E}_T) via which one might be able to search for SUSY. Since the gluino is a Majorana particle, it decays with equal likelihood to positive or negative charginos: the leptonic decays of the chargino can then lead to events with two, isolated, same-sign (SS) charged leptons ^{76,77,78} together with jets plus \cancel{E}_T . If one of the charginos is replaced by a leptonically decaying neutralino, trilepton event topologies result. While other topologies will also be present, the SS and 3ℓ events are especially interesting because (after suitable cuts ^{79,80}) the SM *physics* backgrounds ^z are estimated to be 2.7 fb and 0.7 fb , respectively, for $m_t = 175$ GeV. The corresponding signal cross sections, together with the cross sections in other channels, are illustrated in Fig. 4 which

^zIn addition, there are always detector-dependent instrumental backgrounds from misidentification of jets or isolated pions as leptons, mismeasurement of the sign of the lepton charge *etc.* that a real experimentalist has to contend with.

has been obtained using ISAJET 7.13. It includes signals from all sparticle sources, not just gluinos and squarks. We see that while the cross sections in the clean 3ℓ and SS event topologies are indeed tiny, Tevatron experiments should just about be reaching the sensitivity to probe SUSY via these channels.

Charginos and Neutralinos: The electroweak production of charginos and neutralinos, we have seen, offers yet another channel for probing supersymmetry, the most promising of which is the hadron-free trilepton signal from the reaction $p\bar{p} \rightarrow \tilde{W}_1 \tilde{Z}_2 X$, where both the chargino and the neutralino decays leptonically. In fact, we saw in Fig. 1 that for very large integrated luminosities, this channel potentially offers the maximal reach for supersymmetry (since the OS dilepton signal from $\tilde{W}_1 \tilde{W}_1$ production suffers from large SM backgrounds from WW production). It was first emphasized by Arnowitt and Nath⁸¹ that, with an integrated luminosity of $\sim 100 \text{ pb}^{-1}$, this signal would be observable at the Tevatron even if resonance production of $\tilde{W}_1 \tilde{Z}_2$ is suppressed. A subsequent analyses³⁴ showed that the signal may even be further enhanced in some regions of parameter space due to enhancements in the \tilde{Z}_2 leptonic branching fractions, as discussed in Sec. 5. Detailed Monte Carlo studies^{82,75,35,36} including effects of experimental cuts were performed to confirm that Tevatron experiments should indeed be able to probe charginos via this channel. Indeed an early analysis⁸³ by the CDF collaboration, from a non-observation of this signal, has obtained a limit on the chargino mass that essentially coincides with the one from LEP. While this analysis does not yet lead to an improved bound, it nonetheless shows that Tevatron experiments will eventually probe regions not accessible at LEP, and perhaps, even at LEP2.

This signal, which depends on the neutralino branching fractions, is sensitive to the model parameters, and it is not possible to simply state the reach in terms of the mass of the chargino. For favourable values of parameters, experiments at the MI should be able to probe charginos heavier than 100 GeV, corresponding to $m_{\tilde{g}} \gtrsim 300 - 350 \text{ GeV}$ (at TeV33, up to $\sim 500\text{-}600 \text{ GeV}$ where the two-body spoiler decays of \tilde{Z}_2 become accessible, and for light sleptons, even up to 600-700 GeV); on the other hand, there are other regions of parameter space where the leptonic branching fraction of the neutralino is strongly suppressed^{35,36}, and there is no signal for charginos as light as 45-50 GeV even at the TeV33 upgrade of the Tevatron. Thus, while this channel can probe significant regions of the parameter space of either the MSSM or SUGRA, the absence of any signal in this channel will not allow one to infer a lower limit on $m_{\tilde{W}_1}$.

Top Squarks: We have seen that \tilde{t}_1 , the lighter of the two top squarks, may be rather light so that it may be pair-produced at the Tevatron even if other

squarks and gluinos are all too heavy. If its decay to chargino is allowed, the leptonic decay of one, or both, stops lead to events with one, or two, isolated leptons together with jets plus \cancel{E}_T , exactly the same event topologies as for the top quark search. Thus $t\bar{t}$ production is the major background⁸⁴ to the t -squark search. Because $\sigma(t\bar{t}) \sim 10\sigma(\tilde{t}_1\bar{\tilde{t}}_1)$ for a top and stop of the same mass, stop signals are detectable at the Tevatron only if the stop is considerably lighter than the top. In the single lepton channel it has been shown^{85,86} that with an integrated luminosity of around 100 pb^{-1} , the stop signal should be detectable at the Tevatron if $m_{\tilde{t}_1} \lesssim 100\text{ GeV}$, provided b -jets can be adequately tagged. In the dilepton channel, the t -squark signal can be separated from the top background by searching for *soft* dilepton events: for instance, eliminating events with $|p_T(\ell_1)| + |p_T(\ell_2)| + |\cancel{E}_T| > 100\text{ GeV}$ effectively removes the top background if $m_t > 150\text{ GeV}$, allowing⁸⁵ the search for stops lighter than about 80-100 GeV without the need for b -tagging.

If the chargino is heavy, the stop will instead decay via $\tilde{t}_1 \rightarrow c\tilde{Z}_1$ and stop pair production will be signalled by dijet plus \cancel{E}_T events, and hence looks like the squark signal, but without any cascade decays. Again suitable cuts will allow⁸⁵ Tevatron experiments to probe $m_{\tilde{t}_1} \lesssim 100\text{ GeV}$ with 100 pb^{-1} of data, even if the LSP is relatively heavy. In fact, there is already a preliminary analysis by the D0 collaboration⁸⁷ that excludes $60\text{ GeV} < m_{\tilde{t}_1} < 100\text{ GeV}$ if the LSP mass is smaller than 25-50 GeV.

At the MI, stop masses up to 120-130 GeV should be explorable using essentially the same strategies. Mrenna *et. al.*³⁶ using cuts optimized to detect heavier stops, find that, with an integrated luminosity of 2 fb^{-1} , it should be possible to explore stops as heavy as 160 GeV if they decay via the chargino mode. They claim a reach of 200 GeV, with a data sample of 25 fb^{-1} , that may be available at TeV33. If chargino is too heavy for the decay $\tilde{t}_1 \rightarrow b\tilde{W}_1$ to be accessible but the stop mass is in the 180-250 GeV range, we note that the three body decay $\tilde{t}_1 \rightarrow bW\tilde{Z}_1$ may be kinematically accessible. In this case, one has to see how this compares to the loop decay $\tilde{t}_1 \rightarrow c\tilde{Z}_1$ in order to assess the viability of this signal.

Sleptons: The best hope for slepton detection appears⁸⁸ to be via the clean OS dilepton plus \cancel{E}_T channel. But even here, there is a large irreducible background from WW production as well as possible contamination of the signal from other SUSY sources such as chargino pair production. It was concluded that at the MI it would be very difficult to see sleptons from off-shell Z production; *i.e.* if $m_{\tilde{\ell}} \gtrsim 50\text{ GeV}$. Experiments at TeV33 should probe sleptons with masses up to about 100 GeV, given an integrated luminosity of 25 fb^{-1} .

SUSY Searches at the LHC

While it is certainly possible that SUSY may be discovered at an upgraded Tevatron, there are parameter ranges where a SUSY signal may evade detection even if sparticles are not very heavy. There is no doubt that in order to cover the complete parameter-space of weak scale SUSY either a linear collider operating at a centre of mass energy $\sim 0.5 - 1$ TeV or the LHC is necessary.

We see from Fig. 2 that at the LHC, squarks and gluinos dominate sparticle production for gluino masses beyond 1 TeV: it is thus reasonable to focus most attention on these although, of course, signals should be looked for in all possible channels. For reasons of brevity, and because the ideas involved in LHC searches are qualitatively similar to those described above, we will content ourselves with just presenting an overview of the LHC reach, and refer the interested reader to the vast amount of literature that already exists for details.

As before, the cascade decays of gluinos and squarks result in n -jet plus m -leptons plus \cancel{E}_T events^{89,90} where $m = 0$ corresponds to the classic \cancel{E}_T signal. Of the multilepton channels, the SS and $m \geq 3$ channels suffer the least from SM backgrounds. It is worth keeping in mind that at the LHC, many different sparticle chains contribute to a particular event topology, and further, that the dominant production mechanism for any particular channel depends on the model parameters. For instance, gluino pair production (with each of the gluinos decaying to a chargino of the same sign) are generally regarded as the major source of SS dilepton events; notice, however, that the reaction $pp \rightarrow \tilde{b}_L \tilde{b}_L^* X \rightarrow t \tilde{W}_1^- \bar{t} \tilde{W}_1^+ X$ may also be a copious source of such events, since now the leptons can come either from top or chargino decays (recall that we had noted that \tilde{b}_L may be relatively light). It is, therefore, necessary to simultaneously generate all possible sparticle processes in order to realistically simulate a signal in any particular event topology. This is possible using ISAJET. However, this raises another issue which is especially important at the LHC. If we see a signal in any particular channel, can we uncover its origin? We will return to this later, but for now, focus ourselves on the SUSY reach of the LHC.

The ATLAS collaboration⁹⁰ at the LHC has done a detailed analysis of the signal in the \cancel{E}_T as well as in the SS dilepton channels. They found that gluinos as light as 300 GeV should be easily detectable in the \cancel{E}_T channel. Then requiring rather stiff cuts, $\cancel{E}_T > 600$ GeV, $p_T(\text{jet}_1, \text{jet}_2, \text{jet}_3) > 200$ GeV, $p_T(\text{jet}_4) > 100$ GeV along with a cut $S_T > 0.2$ on the transverse sphericity, they find that it should be possible to search for gluinos with a mass up to 1.3 TeV (2 TeV) for $m_{\tilde{q}} = 2m_{\tilde{g}}$ ($m_{\tilde{q}} = m_{\tilde{g}}$), assuming an integrated luminosity of 10 fb^{-1} . This reach changes is altered by about ± 300 GeV if the integrated luminosity is changed by an order of magnitude. Very similar results for the

reach in the \cancel{E}_T channel have also been obtained⁹¹ within the context of the SUGRA framework, although the event selection criteria used are quite different. In the same-sign dilepton channel, the ATLAS collaboration⁹⁰, concludes that the reach of the LHC will be 900-1400 GeV (for $m_{\tilde{q}} = 2m_{\tilde{g}}$) or 1100-1800 GeV (for $m_{\tilde{q}} = m_{\tilde{g}}$), where the lower (higher) number corresponds to a luminosity of 1 fb^{-1} (100 fb^{-1}). Prospects for SUSY detection in the SS and other multilepton channels (with or without real Z bosons) have also been discussed⁸⁹ by other authors. An analysis of the multilepton signals within the context of SUGRA models is in progress. Preliminary results from this analysis⁹² indicate that the SUSY reach in the 1ℓ channel may extend beyond that in the \cancel{E}_T channel. While this may appear somewhat surprising at first sight because this channel is plagued by large backgrounds from $W \rightarrow \ell\nu$ and $t\bar{t}$ production, these authors have exploited the presence of hard jets and large \cancel{E}_T in SUSY events to devise cuts that reduce this background to a manageable level.

Within the SUGRA framework, the LHC should, in the clean trilepton channel, be able^{93,94} to probe $\widetilde{W}_1\widetilde{Z}_2$ production all the way up to where spoiler modes of \widetilde{Z}_2 become accessible if $\mu < 0$ and $\tan\beta$ is not too large. Then, it is possible to find a set of cuts that cleanly separate the $\widetilde{W}_1\widetilde{Z}_2$ event sample from both SM backgrounds as well as other sources of SUSY events. This will prove important later. For positive values of μ , signals are readily observable for rather small and large values of m_0 ; in the intermediate range $400 \text{ GeV} \lesssim m_0 \lesssim 1000 \text{ GeV}$, this signal is suppressed because of the suppression of the leptonic \widetilde{Z}_2 branching fraction emphasized earlier.

As at Tevatron upgrades, the OS dilepton channel offers^{95,88} the best opportunity for slepton searches. At the LHC, it should be possible to detect sleptons up to about 250-300 GeV, although excellent jet vetoing capability will be needed to detect the signal for the highest masses.

The SUSY reach at various possible future facilities is summarized in the Table.

Several comments are worth noting:

- In some places, two sets of numbers are given for the reach. These correspond to results from different analyses more fully described in the review⁴ from where this Table is taken. Basically, the more conservative number also requires the signal to be larger than 25% of the background, in addition to exceeding the 5σ level. Also, the two analyses do not use the same cuts.
- The multilepton rates in the Table are shown for negative values of μ and $\tan\beta = 2$. For other parameters, especially for $\mu > 0$, the trilepton

Table 1: Estimates of the discovery reach of various options of future hadron colliders. The signals have mainly been computed for negative values of μ . We expect that the reach in especially the $all \rightarrow 3\ell$ channel will be sensitive to the sign of μ .

Signal	Tevatron II 0.1 fb ⁻¹ 1.8 TeV	Main Injector 1 fb ⁻¹ 2 TeV	TeV33 10 fb ⁻¹ 2 TeV	LHC 10 fb ⁻¹ 14 TeV
$\cancel{E}_T(\tilde{q} \gg \tilde{g})$	$\tilde{g}(210)/\tilde{g}(185)$	$\tilde{g}(270)/\tilde{g}(200)$	$\tilde{g}(340)/\tilde{g}(200)$	$\tilde{g}(1300)$
$l^\pm l^\pm(\tilde{q} \gg \tilde{g})$	$\tilde{g}(160)$	$\tilde{g}(210)$	$\tilde{g}(270)$	
$all \rightarrow 3l(\tilde{q} \gg \tilde{g})$	$\tilde{g}(180)$	$\tilde{g}(260)$	$\tilde{g}(430)$	
$\cancel{E}_T(\tilde{q} \sim \tilde{g})$	$\tilde{g}(300)/\tilde{g}(245)$	$\tilde{g}(350)/\tilde{g}(265)$	$\tilde{g}(400)/\tilde{g}(265)$	$\tilde{g}(2000)$
$l^\pm l^\pm(\tilde{q} \sim \tilde{g})$	$\tilde{g}(180 - 230)$	$\tilde{g}(320 - 325)$	$\tilde{g}(385 - 405)$	$\tilde{g}(1000)$
$all \rightarrow 3l(\tilde{q} \sim \tilde{g})$	$\tilde{g}(240 - 290)$	$\tilde{g}(425 - 440)$	$\tilde{g}(550)$	$\gtrsim \tilde{g}(1000)$
$\tilde{t}_1 \rightarrow c\tilde{Z}_1$	$\tilde{t}_1(80-100)$	$\tilde{t}_1(120)$	$\tilde{t}_1(150)$	
$\tilde{t}_1 \rightarrow b\tilde{W}_1$	$\tilde{t}_1(80 - 100)$	$\tilde{t}_1(120)$	$\tilde{t}_1(180)$	
$\Theta(\tilde{t}_1\tilde{t}_1^*) \rightarrow \gamma\gamma$	—	—	—	$\tilde{t}_1(250)$
$4\tilde{l}\tilde{l}^*$	$\tilde{l}(50)$	$\tilde{l}(50)$	$\tilde{l}(100)$	$\tilde{l}(250-300)$

rates may be strongly suppressed due to a suppression of the \tilde{Z}_2 branching fraction discussed above. Notice also that at TeV33, the reach in the leptonic channels exceeds that in the \cancel{E}_T channel. At the TeV33 upgrade, hadronically quiet trilepton events may be observable all the way up to the spoiler modes for favourable ranges of model parameters. It is, however, important to remember that supersymmetry may escape detection at these facilities even if sparticles are relatively light.

- At the LHC, gluinos and squarks are detectable to well beyond 1 TeV in the \cancel{E}_T channel, and up to 2 TeV if their masses are roughly equal. Thus the LHC should be able to probe the complete parameter space of weak scale SUSY, at least within the assumed framework. Moreover, there should be some observable signals in the leptonic channels if a signal in the \cancel{E}_T channel is to be attributed to supersymmetry.
- Tevatron upgrades will not probe sleptons significantly beyond the reach of LEP2, whereas the LHC reach may be comparable to that of the initial phase of linear colliders.
- Tevatron upgrades should be able to detect \tilde{t}_1 with a mass up to 120 GeV at the MI⁸⁵, and up to 150-180 GeV at TeV33.³⁶ It has also been pointed out, assuming that $\tilde{t}_1 \rightarrow c\tilde{Z}_1$ is its dominant decay, that it should be possible⁹⁶ to search for \tilde{t}_1 at the LHC via the two photon decay of the scalar $\tilde{t}_1\tilde{t}_1^*$ bound state, in much the same way that Higgs bosons searches (to be discussed next) are carried out.

Higgs Bosons: At the LHC, MSSM Higgs bosons are dominantly produced¹⁰ by gg fusion (via loops of quarks and squarks), and for some parameter ranges, also via $b\bar{b}$ fusion. Vector boson fusion, which in the SM dominates these other mechanisms for large Higgs masses ($m \gtrsim 600$ GeV) is generally unimportant, since the couplings of heavy Higgs bosons to VV pairs is suppressed. Unfortunately, we do not have much time to discuss various strategies that have been suggested⁹⁷ for the detection of the Higgs sector of SUSY. Over much of the parameter space, all the Higgs bosons except the lightest neutral scalar, H_ℓ , are too heavy to be of interest, although for some ranges of MSSM parameters, signals from $H_h \rightarrow \gamma\gamma, \tau\bar{\tau}, \mu\bar{\mu}, 4\ell$ and $H_p \rightarrow \gamma\gamma, \tau\bar{\tau}, \mu\bar{\mu}$ may be observable. The two photon decay mode is the most promising channel for H_ℓ detection at the LHC. The region of parameter space where there is some signal for an MSSM Higgs boson either at the LHC or at LEP2 have been nicely summarized in the technical report of the CMS Collaboration⁹⁸, assuming that sparticles are too heavy to be produced via Higgs boson decays.^a The most striking feature of their analysis is that despite optimistic detector assumptions, there are regions of parameter space where *there may be no signal for any of the Higgs bosons either at the LHC or at LEP2*. Part of this hole may be excluded⁶³ by analyses of rare decays such as $b \rightarrow s\gamma$ mentioned earlier. It has been suggested¹⁰¹ that Higgs boson signals may also be detectable in this hole region via $t\bar{t}H_\ell$ production where a lepton from t decay may be used to tag the event so that H_ℓ can then be detected via its dominant $b\bar{b}$ decay. This would require efficient b tagging with the high luminosity option for the LHC. Whether this is technically possible is not clear at this time. It is worth remembering that Higgs boson detection would be relatively easy at a linear collider, the first of many examples of the complementary nature of these facilities.

10 Beyond SUSY Discovery: More Ambitious Measurements

We have seen that if the minimal SUSY framework that we have adopted is a reasonable approximation to nature, experiments at supercolliders should certainly be able to detect signals for physics beyond the SM. If we are lucky, such signals might even show up at LEP2 or at Tevatron upgrades. We will then have to figure out the origin of these signals. If the new physics is super-

^aThis is not necessarily a good assumption. The branching fractions for the SUSY decays can be quite substantial for large regions of parameter space, and can reduce the signals via which the Higgs bosons are usually searched for and increase the parameter space hole referred to below.⁹⁹ Sometimes, however, they lead to novel signals for Higgs boson searches which can then refill¹⁰⁰ some of the hole region. Of course, for almost all cases where SUSY Higgs decays are important, it should also be possible to detect the sparticles at the LHC. It is only Higgs boson detection that may be more difficult.

symmetry, it is likely (certain, at the LHC) that there will simultaneously be signals in several channels. While the observation of just one or two of these signals would convince the believers, others would probably demand stronger evidence. It is not, however, reasonable to expect that we will immediately detect all (or even several of) the super-partners. Thus, it is important to think about just what information can be obtained in various experiments, information that will help us to elucidate the nature of the underlying physics. Towards this end, we would like to be able to,

- measure any new particle's masses and spins, and
- measure its couplings to SM particles; these would serve to pin down its internal quantum numbers.

More ambitiously, we may ask:

- Assuming that the minimal framework we have been using is correct, is it possible to measure the model parameters? Is it possible to actually provide tests for, say, the minimal SUGRA framework, and thus also test the assumptions about the physics at the GUT or Planck scale that are an integral part of this picture?
- At hadron colliders, especially, where several new sparticle production mechanisms may be simultaneously present,^b is it possible to untangle these from one another?
- As mentioned in Sec. 1, like any other (spontaneously broken) symmetry, supersymmetry, though softly broken, implies relationships between the various couplings in the theory. Is it possible to directly test supersymmetry by experimentally verifying these coupling constant relationships?

10.1 Mass Measurements

e^+e^- colliders: The clean environment of e^+e^- colliders as well as the very precise energy of the beam allows for measurements of sparticle masses. We will briefly illustrate the underlying ideas with a simple example. It is easy to show that the total cross section for smuon production has the energy dependence,

$$\sigma(\tilde{\mu}\tilde{\mu}^*) \propto \left(1 - \frac{4m_{\tilde{\mu}}^2}{s}\right)^{\frac{3}{2}},$$

^bSince the energy of the linear collider is likely to be increased in several steps to the TeV scale, one may hope that this will be less of a problem there. The lighter sparticles will be discovered first. Knowledge about their properties thus obtained should facilitate the untangling of the more complex decays of heavier sparticles.

and further, that the energy distribution of the daughter muon from the decay of the smuon, assuming only direct decays to the LSP, is flat and bounded by¹⁰²

$$\frac{m_{\tilde{\mu}}^2 - m_{\tilde{Z}_1}^2}{2(E + p)} \leq E_{\mu} \leq \frac{m_{\tilde{\mu}}^2 - m_{\tilde{Z}_1}^2}{2(E - p)},$$

with $E(p)$ being the energy (momentum) of the smuon. We thus see that the energy dependence of the smuon cross section gives a measure of the smuon mass, while a measurement of the end points of the muon energy spectrum yields information about $m_{\tilde{\mu}}$ as well as $m_{\tilde{Z}_1}$. Of course, theoretically these relations are valid for energy and momentum measurements made with ideal detectors without any holes and with perfect energy and momentum resolutions. In real detectors there would be smearing effects as well as statistical fluctuations. It has been shown,⁶¹ taking these effects into account, that with an integrated luminosity of 100 pb^{-1} , experiments at LEP2 should be able to determine the smuon mass within 2-3 GeV. At the JLC, an integrated luminosity of 20 fb^{-1} should enable^{51,72} the determination of the smuon and LSP masses to within 1-2 GeV. If sleptons are heavy but charginos light, a study of the reaction $e^+e^- \rightarrow \widetilde{W}_1 \widetilde{W}_1 \rightarrow jj\tilde{Z}_1 + \ell\nu\tilde{Z}_1$ should allow the determination of $m_{\tilde{W}_1}$ and $m_{\tilde{Z}_1}$ with a precision of $\lesssim 3 \text{ GeV}$, both at LEP2 (where an integrated luminosity of about 1 fb^{-1} would be necessary) and at the JLC. A good jet mass resolution is crucial. Finally, it has also been shown¹⁰³ that with the availability of beam polarization at linear colliders it should be possible to determine squark masses with a precision of $\sim 5 \text{ GeV}$ even if these decay via MSSM cascades: in particular, it should be possible to determine the splittings amongst the squarks with good precision.

Hadron Colliders: Can one say anything about sparticle masses from SUSY signals at hadron colliders? Despite the rather messy environment, it is intuitively clear that this should be possible if one can isolate a single source of SUSY events from both SM backgrounds as well as from other SUSY sources: a study of the kinematics would then yield a measure of sparticle masses within errors determined by the detector resolution. We have already seen that it is indeed possible to isolate a relatively clean sample of $p\bar{p} \rightarrow \widetilde{W}_1 \tilde{Z}_2 \rightarrow \ell\bar{\ell}\ell' + \cancel{E}_T$ events at the LHC. The end-point of the $m_{\ell\bar{\ell}}$ distribution, it has been shown,^{93,94} yields an accurate measure of $m_{\tilde{Z}_2} - m_{\tilde{Z}_1}$. It may also be possible to extract relationships between the masses of charginos and neutralinos, though this is less straightforward.

A measurement of the gluino mass would be especially important at hadron colliders since gluinos cannot be pair produced by tree-level processes at e^+e^- colliders. The first attempt⁴⁴ involved a parton-level examination of events

from $\tilde{g}\tilde{Z}_1$ production, with a very hard \cancel{E}_T cut to separate these events from $\tilde{g}\tilde{g}$ events. It was argued that the mass of the hadronic system recoiling against the missing transverse energy should yield $m_{\tilde{g}}$. The analysis suffered from the fact that QCD radiation as well as other SUSY processes which could contaminate the $\tilde{g}\tilde{Z}_1$ sample were not included. It was ultimately concluded that this reaction might be of use but only if $m_{\tilde{g}} < 350$ GeV. Another strategy⁷⁷ focussed on the isolated same sign dilepton sample from gluino pair production (which is expected to be relatively free of SM backgrounds). Each event was divided into two hemispheres defined by the transverse sphericity axis, and an estimator of $m_{\tilde{g}}$ constructed. It was claimed that $m_{\tilde{g}}$ could be measured with a precision of about 15%. This conclusion may be overly optimistic since this study considered just a single SUSY source of same sign dileptons (we have seen several sources of such events), a single cascade decay chain ($\tilde{g} \rightarrow q\bar{q}\tilde{W}_1$, $\tilde{W}_1 \rightarrow \ell\nu\tilde{Z}_1$) and neglected effects of QCD radiation. This hemispheric separation strategy was adopted in a recent attempt⁹¹ to obtain $m_{\tilde{g}}$ from the \cancel{E}_T event sample. The larger of the masses of the hadronic system in the two hemispheres where events were selected requiring at least two jets and \cancel{E}_T larger than a preassigned value E_T^c was used as an estimator for $m_{\tilde{g}}$. All sources of SUSY events and QCD radiation were simulated using ISAJET. It was claimed that the gluino mass may be measured with a precision of 15-25% provided $m_{\tilde{g}} \lesssim 800$ GeV, beyond which the cross section becomes too small to be able to reconstruct distributions with sufficient precision.

10.2 Determination of Spin at e^+e^- colliders

If sparticle production occurs via the exchange of a vector boson in the s -channel, it is easy to check that the sparticle angular distribution is given by, $\sin^2\theta$ for spin zero particles, and $E^2(1 + \cos^2\theta) + m^2\sin^2\theta$ for spin half particles. Thus if sparticles are produced with sufficient boost, the angular distribution of their daughters which will be relatively strongly correlated to that of the parent, should be sufficient to distinguish between the two cases. Chen *et. al*⁶¹ have shown that, with an integrated luminosity of 500 pb^{-1} , it should be possible to determine the smuon spin at LEP2. A similar analysis⁷² has been performed for a 500 GeV linear collider.

10.3 Testing the Minimal SUGRA GUT Framework

We begin by recalling that all the minimal SUGRA GUT framework is determined by just the four parameters, m_0 , $m_{\frac{1}{2}}$, $\tan\beta$ and A_0 which, together with the sign of μ completely determine all the sparticle masses and couplings.

Since the number of observables can be much larger than the number of parameters, there must exist relations between observables which can be subjected to experimental tests. In practice, such tests are complicated by the fact that there are experimental errors, and further, it may not be possible to cleanly separate between what, in principle, should be distinct observables; *e.g.* cross sections for E_T events from $\tilde{g}\tilde{g}$, $\tilde{g}\tilde{q}$ and $\tilde{q}\tilde{q}$ sources at the LHC.

Because of the clean experimental environment and the availability of polarized beams, these tests can best be done at e^+e^- colliders, where we have already seen that it is possible to determine sparticle masses with a precision of 1-2%. The determination of the selectron and smuon masses will allow us to test their equality $m_{\tilde{e}_L} = m_{\tilde{\mu}_L}$, $m_{\tilde{e}_R} = m_{\tilde{\mu}_R}$ at the percent level^{51,72} — the same may be done with staus, though with a somewhat smaller precision. This is a test of the assumed universality of slepton masses.

A different test may be possible if both $\tilde{\ell}_R$ and \tilde{W}_1 are kinematically accessible and a right-handed electron beam is available. It is then possible to measure $m_{\tilde{Z}_1}$, $m_{\tilde{W}_1}$, $\sigma_R(\tilde{\ell}_R\tilde{\ell}_R)$ and $\sigma_R(\tilde{W}_1\tilde{W}_1)$ (note that the chargino cross section for right-handed electron beams has no contribution from sneutrino exchange!). These four observables can then be fitted to the four MSSM parameters μ , $\tan\beta$ and the electroweak gaugino masses μ_1 and μ_2 . In practice^{51,72}, while μ may be rather poorly determined if the chargino is dominantly a gaugino, $\frac{\mu_1}{\mu_2}$ is rather precisely obtained so that it should be possible to test the gaugino mass unification condition at the few percent level, given 50 fb^{-1} of integrated luminosity.^c It may further be possible to determine $m_{\tilde{\nu}_e}$ by measuring chargino production with a left-handed electron beam. This is of interest because the difference between the squared masses of the sneutrino and $\tilde{\ell}_L$ is a direct test of the SU(2) gauge symmetry for sleptons. For further details and other interesting tests, we refer the reader to the original literature.^{51,72}

Analogous tests are much more difficult at hadron colliders. Nevertheless it may be possible to combine the data from the Tevatron and LEP experiments (again note the complementarity between hadron and e^+e^- colliders) to test for consistency of the SUGRA framework. Since the signals are fixed in terms of just four parameters (and a sign) it is convenient¹⁰⁴ to display these in the $m_0 - m_{\frac{1}{2}}$ plane for fixed values of $\tan\beta$ and A_0 . That way, various correlations become obvious. For example, it is easy to find regions where several signals should be simultaneously present, and at roughly predicted levels. Observation (or non-observation) of such signals at various facilities would serve to test these correlations and also, perhaps, begin to zero in on the model parameters.

^cFeng and Strassler⁴⁹ have shown that, with 1 fb^{-1} of data, a test of this relation at the 20% level may also be possible at LEP2.

Of course, a determination of these is a complex and difficult task, and only very preliminary work has been done on this issue.

10.4 Identifying Sparticle Production Mechanisms at the LHC

At e^+e^- colliders where the centre of mass energy is incrementally increased, it may be reasonable to suppose that it is unlikely (except, perhaps, for the sfermion degeneracy expected in SUGRA type models) that several particle thresholds will be crossed at the same time. It would thus be possible to focus on just one new signal at a time, understand it and then proceed to the next stage. The situation at the LHC will, of course, be quite different. Several sparticle production mechanisms will be simultaneously present as soon as the machine turns on, so that even if it is possible to distinguish new physics from the SM, the issue of untangling the various sparticle production mechanisms will remain. For example, even if we attribute a signal in the $\cancel{E}_T + jets$ channel to sparticle production, is it possible to tell whether the underlying mechanism is the production of just gluinos or of gluinos and squarks? ^d

Some progress has already been made in this direction. We have already seen that the $\tilde{W}_1\tilde{Z}_2$ source of trileptons can clearly be isolated from other SUSY processes. The opposite sign dilepton signal from slepton production is probably distinguishable from the corresponding signal from chargino pair production since the dileptons from slepton production always have the same flavour. ^e To tell whether squarks are being produced in addition to gluinos, at least two distinct strategies have been suggested. The first makes use of the fact that there are more up quarks in the proton than down quarks. We thus expect many more $\tilde{u}_L\tilde{u}_L$ and $\tilde{g}\tilde{u}_L$ events as compared to $\tilde{d}_L\tilde{d}_L$ and $\tilde{d}_L\tilde{g}$ events at the LHC. As a result, any substantial production of squarks in addition to gluinos will be signalled ¹⁰⁵ by a charge asymmetry in the same-sign dilepton sample: cascade decays of gluinos and squarks from $\tilde{g}\tilde{q}$ and $\tilde{q}\tilde{q}$ events lead to a larger cross section for positively charged same sign dileptons than for negatively charged ones. This has since been confirmed by detailed studies by the ATLAS collaboration ⁹⁰ where the SS dilepton charge asymmetry is studied as a function of $\frac{m_{\tilde{q}}}{m_{\tilde{g}}}$, and shown to monotonically disappear as this ratio becomes small. More recently, it has been suggested ⁹¹ that a study of the jet multiplicity in the \cancel{E}_T sample could also reveal the production of squarks provided one has some idea about the gluino mass. The idea is to note that \tilde{q}_R , which

^dHere, we tacitly assume that squarks will not be much lighter than gluinos.

^eThe extent to which this channel is contaminated by other SUSY sources has not been explicitly checked. The viability of the opposite sign, clean dilepton signal from chargino pair production at the LHC is under investigation.

are produced as abundantly as \tilde{q}_L , frequently decay directly to the LSP via $\tilde{q}_R \rightarrow q\tilde{Z}_1$ and so lead to only one jet (aside from QCD radiation). In contrast, gluinos decay via $\tilde{g} \rightarrow q\tilde{q}\tilde{W}_i$ or $\tilde{g} \rightarrow q\tilde{q}\tilde{Z}_i$, so that that gluino decays contain two, and frequently more, jets from their cascade decays. Thus the expected jet multiplicity is lower if squark production forms a substantial fraction of the \cancel{E}_T sample. Of course, since $\langle n_{jet} \rangle$ (from gluino production) depends on its mass, some idea of $m_{\tilde{g}}$ is necessary for this strategy to prove useful. A detailed simulation⁹¹ shows that the mean value of the n_{jet} distribution increases by about $\frac{1}{2}$ unit, when the squark mass is increased from $m_{\tilde{q}} = m_{\tilde{g}}$ by about 60-80%.

Cascade decays of gluinos and squarks could also lead to the production of the Higgs bosons of supersymmetry. It is, therefore, interesting to ask whether these can be detected in the data sample which has already been enriched in SUSY events. Neutral Higgs bosons might be detectable¹⁰⁵ via an enhancement of the multiplicity of central b -jets in the \cancel{E}_T or same sign dilepton SUSY samples. Some care must be exercised in drawing conclusions from this because such enhancements may also result because third generation squarks happen to be lighter than the other squarks.^f It has also been shown⁹¹ that it may also be possible to reconstruct a mass bump in the $m_{b\bar{b}}$ distribution if there is a significant branching fraction for the decay $\tilde{Z}_2 \rightarrow H_\ell \tilde{Z}_1$ and H_ℓ is produced in events with no other b -jets since then we would have a large combinatorial background. The charged Higgs boson, if it is light enough, may be identifiable via the detection of τ lepton enhancements in SUSY events¹⁰⁵; it should, however, be kept in mind that such light charged Higgs bosons also contribute to the $b \rightarrow s\gamma$ decays.

We stress that the complex cascade decay chains of gluinos and squarks may be easier to disentangle if we already have some knowledge about the masses and couplings of the lighter charginos and neutralinos that are produced in these decays. While it is indeed possible that \tilde{W}_1 may be discovered at LEP2 and that its mass is determined there, it is likely that we may have to wait for experiments at the linear collider to be able to pin down the couplings, and, perhaps, even for discovery of \tilde{W}_1 and \tilde{Z}_2 . In this case, a reanalysis of the LHC data in light of new information that may be gained from these experiments may prove to be very worthwhile: it may thus be necessary to archive this data in a form suitable for subsequent reanalysis. Once again, we see the complementary capabilities of e^+e^- and hadron colliders.

^fThese may be directly produced with large cross sections or may lead to enhancement of gluino decays to third generation fermions as discussed in Sec. 5.

10.5 Direct Tests of Supersymmetry

We have already seen that supersymmetry, like any other symmetry, implies relationships⁹ between various dimensionless couplings in the theory even if it is softly broken. For example, the fermion-sfermion-gaugino (or, since the Higgs multiplet also forms a chiral superfield, the Higgs-Higgsino-gaugino) coupling is completely determined by the corresponding gauge coupling. A verification of the relation would be a direct test of the underlying supersymmetry. We emphasize that such a test would be essentially model independent as it relies only on the underlying global supersymmetry, and not on any details such as assumptions about physics at the high scale or even the sparticle content. The main complication is that the gauginos (or the Higgs bosons and Higgsinos, or for that matter, the sfermions) are not mass eigenstates, so that the mixing pattern has to be disentangled before this test can be applied. This will require an accurate measurement of several observables which can then be used to disentangle the mixing and also simultaneously to measure the relevant coupling.

Feng *et. al.*¹⁰⁷ have argued that such a test is best done via a determination of chargino properties. As we have seen, the charged gaugino and the corresponding Higgsino can mix only if electroweak symmetry is broken. This is the reason why the off-diagonal terms in the chargino matrix are equal to⁶ $\sqrt{2}M_W \cos \beta$ and $\sqrt{2}M_W \sin \beta$, respectively. Assuming that the chargino is a mixture of just one Dirac gaugino and one Dirac Higgsino, the most general mass matrix would contain four parameters: the two diagonal elements and the two off-diagonal ones. These latter can always be parametrized by $\sqrt{2}M_W^\chi \cos \beta^\chi$ and $\sqrt{2}M_W^\chi \sin \beta^\chi$. It is the SUSY constraint on the Higgs-Higgsino-gaugino coupling that forces $M_W^\chi = M_W$. Within the MSSM framework (which we adopt for evaluating the feasibility of the SUSY test), for parameters such that the chargino is a substantial mixture of the gaugino and Higgsino, both charginos should be accessible at a 500 GeV linear collider. A determination of four quantities, chosen to be the masses of the two charginos along with the total production cross section from right-handed electron beam (recall this does not couple to the sneutrino, so that the cross section is determined by gauge interactions) and an appropriately defined¹⁰⁷ forward-backward asymmetry is thus sufficient to determine the four entries of the chargino mass matrix. With an integrated luminosity of 30 fb^{-1} , the relation $M_W^\chi = M_W$ can be tested at the 30% level at a 500 GeV linear collider. We will refer the reader to the original paper¹⁰⁷ for further details and also an

⁹These relations are corrected by radiative corrections which are generally expected to be smaller than a few percent.¹⁰⁶

analogous test (which can be done also at about the 30% level) for parameters such that the chargino is mainly a gaugino. Finally, if the chargino is Higgsino-like, the lightest neutralino is essentially degenerate with it (within the MSSM framework). In this case, even the observation of the chargino signal (let alone precision tests) may be difficult because the decay products tend to be very soft.

11 Beyond Minimal Models

Up to now, we have confined our analysis to the MSSM framework. Even here, we saw in Sec. 3 that the unmanageably large number of free parameters required us to make additional assumptions in order to obtain tractable phenomenology. It is clearly impractical to seriously discuss the phenomenology of various extensions of the MSSM framework. Here we will merely list some of the ways in which this framework may be modified, and leave it to the reader to figure out the implications for phenomenology. Thinking about this will also help to view our previous discussion in proper perspective.

The MSSM framework may be extended or modified in several ways.

- We may give up the exact universality of the gaugino masses at the GUT scale. Threshold corrections due to unknown GUT, and perhaps even gravitational, interactions would certainly yield model-dependent corrections¹⁰⁸ which preclude exact unification. It is also conceivable that there is, in fact, no grand unification at all, but the observed unification of couplings in LEP experiments is a result¹⁰⁹ of string type unification; in this case, the gaugino masses need not be identical exactly at M_{GUT} . We have already noted that even in SUGRA type models, we do not really know the exact scale at which scalar masses unify. We also remark that the assumption of minimal kinetic energy terms is crucial for obtaining universal scalar masses. Since the Lagrangian for supergravity is non-renormalizable, this really is an additional¹¹⁰ assumption. Any deviation from the assumed universality of scalar masses will, at the very least, modify the conditions of radiative symmetry breaking. Finally, models have also been constructed¹¹¹ where SUSY breaking is a relatively low scale ($\sim 10 - 100$ TeV, as opposed to M_X) phenomenon, so that the scalar mass patterns will be quite different from those in minimal SUGRA.
- R -parity may be explicitly broken by superpotential interactions g_1 and g_2 in Eq. (4) and Eq. (5).

- There could be additional chiral superfields: new generations (with heavy neutrinos), additional Higgs multiplets, or a right-handed sneutrino superfield. We certainly do not need new generations or new Higgs doublets, as they may spoil the observed unification of couplings. Higgs fields in higher representations cause additional problems if they develop a vacuum expectation value. Higgs singlets cannot be logically excluded, and are interesting because they allow for new quartic Higgs boson couplings, though one would have to understand what keeps them from acquiring GUT or Planck scale masses. A singlet right-handed sneutrino (note that this is *not* the superpartner of the usual neutrinos) is an interesting possibility since it occurs in SO(10) GUT models, and also, because it allows for spontaneous breaking of R -parity conservation.¹¹²
- Finally, we could consider models with extended low energy gauge groups — either left-right symmetric models¹¹³ or models with additional Z bosons.¹¹⁴

For want of time and space, we will only qualitatively review how our earlier discussion is altered if R -parity conservation is explicitly violated.¹¹⁵ In some sense, this is the minimal extension because it does not require the introduction of any new particles. Notice, however, that a general analysis of this requires the introduction of 45 new couplings: 9 λ 's, 27 λ' 's and 9 λ'' 's. There are relations amongst these couplings in theories with larger symmetries; *e.g.* GUTs. As we have already discussed, many products of the baryon- and lepton-number violating couplings are strongly constrained. In phenomenological analyses, it is customary (and in light of the large number of new parameters, convenient) to assume that one of the couplings dominates. Even so, several of the couplings are strongly constrained. In a very nice analysis, Barger *et al.*¹¹⁶ have studied the implications from various experiments — β -decay universality, lepton universality, $\nu_\mu e$ scattering, e^+e^- forward-backward asymmetries and ν_μ deep-inelastic scattering — for these new interactions. They find strong constraints on the lepton-number violating couplings, assuming^h that only one of the couplings is non-zero: for instance, they find that of the λ -type couplings, only λ_{131} and λ_{133} can exceed 0.2 (compare this with the electromagnetic coupling $e = 0.3$) for a SUSY scale of 200 GeV, though several λ' and many more of the λ'' interactions can exceed this value. Dimopoulos and Hall¹¹⁷ have, from the upper limit on the mass of ν_e , obtained a strong bound ($< 10^{-3}$) on λ_{133} . The same argument yields significant bounds¹¹⁸ on many of

^hConstraints from $\mu \rightarrow e\gamma$ or $\mu \rightarrow 3e$ decays are much stronger if, say, both e and μ number violating interactions are large.

the λ' couplings so that of these only $\lambda'_{112}, \lambda'_{121}$ and λ'_{111} can exceed 0.2. Generally, only the first generation baryon number violating interactions are strongly constrained from the non-observation of $n - \bar{n}$ oscillations or $NN \rightarrow K\bar{K}X$.¹¹⁹

The reason to worry about all this is that if R -parity is not conserved, both sparticle production cross sections as well as decay patterns may be altered. For instance, if λ' interactions are dominant (with $i = 1$), squarks can be singly produced as resonances in ep collisions at HERA¹²⁰, or in the case of λ'' interactions, at hadron colliders.¹²¹ The production rates will, of course, be sensitive to the unknown R -parity violating couplings. Likewise, if R -parity violating couplings are large compared to gauge couplings, these R -violating interactions will completely alter sparticle decay patterns.

Even if all the λ 's are too small (relative to gauge couplings) to significantly affect the production and decays of sparticles (other than the LSP), these interactions radically alter the phenomenology because the LSP decays visibly, so that the classic \cancel{E}_T signature of SUSY is no longer viable. Clearly, the phenomenology depends on the details of the model. Two extreme cases where the LSP decays either purely leptonically into e 's or μ 's and neutrinos via λ -type couplings, or when it always decays into jets via λ'' couplings have been examined for their impact on Tevatron^{79,122} and LHC¹²³ searches for supersymmetry. The signals, in the former case, are spectacular since the decays of each LSP yields two leptons in addition to any other leptons from direct decays of \tilde{W}_1 or \tilde{Z}_2 produced in the gluino or squark cascade decays. With an integrated luminosity of 100 pb^{-1} that has already been accumulated, experiments at the Tevatron should be able to probe gluinos as heavy as 500-600 GeV. In the other case where the LSP decays purely hadronically, gluino and squark detection is much more difficult than in the MSSM. The reason is that the \cancel{E}_T signal is greatly reduced since neutrinos are now the only sources of \cancel{E}_T . In fact, if squarks are heavy, there may well be no reach in this channel even at the Main Injector. Further, the multilepton signals from cascade decays are also degraded because the jets from LSP decays frequently spoil the lepton isolation. Indeed if squarks are heavy, none of the SUSY signals would be observable in this run of the Tevatron; even the Main Injector will then not probe gluino masses beyond ~ 200 GeV (350 GeV, if $m_{\tilde{q}} = m_{\tilde{g}}$). At the LHC, attention has solely been focussed¹²³ on the same-sign dilepton signal from gluino pair production. In the case where the LSP decays purely leptonically, the gluino mass reach is in excess of 1 TeV; in the less favourable case where the LSP decays hadronically, this signal is again much less promising. Sparticle detection should not be a problem in the clean environment of e^+e^- colliders, even if R -parity is not conserved. In fact, LEP should be able to probe regions of parameters not explorable in the MSSM since signals from

LSP pair production can now be detected.¹²⁴

12 Concluding Remarks

We have seen that experiments at the LHC should be able to explore essentially the whole parameter space of weak scale supersymmetry if we require that sparticles provide the degrees of freedom that stabilize the electroweak symmetry breaking sector. While experiments at Tevatron upgrades (or for that matter even at the current Tevatron or at LEP2) will explore substantial regions of this parameter space, and maybe even discover sparticles, a non-observation of any signal should not be regarded as disheartening: the expected mass scale is several hundred GeV up to a TeV, and so may well not be accessible except at supercolliders. Electron-positron linear colliders, with a centre of mass energy of 500-1000 GeV should also be able to discover sparticles (almost certainly so if the frequently assumed unification condition for gaugino masses is correct). Linear colliders are the ideal facility for the discovery and subsequent detailed study of Higgs bosons.

The SUGRA GUT model that we have described in Sec. 4 provides a very attractive and economic framework. It is consistent with known phenomenology, with grand unification and can incorporate (though not explain) the observed pattern of electroweak symmetry breaking. The simplest such model leads to a degeneracy between the first two squark generations and so is automatically consistent with constraints on FCNC. Furthermore, because SUSY is a decoupling theory in that virtual effects of sparticles become suppressed if their masses are much larger than M_Z , the observed agreement of the SM with LEP constraints is simply incorporated. These models also provide a natural candidate for cold dark matter. We have seen, however, that these rather predictive models are based on several assumptions about the physics at very high scales. It is important to keep in mind that one, or more, of these assumptions may prove to be incorrect. This is especially important when considering the design of future high energy physics facilities. While it is reasonable to use the model as a guide, it is important to examine just how sensitively the various signals depend on these assumptions. The important thing, however, is that these assumptions will be testable in future experiments. It is here that the complementary capabilities of hadron colliders and electron-positron linear colliders play a crucial role. The experimental verification of any of these assumptions will provide a window to the symmetries of physics at ultra-high energy scales.

Lest the preceding discussion gives an impression that SUSY solves most of the problems of particle physics, we should remember that it addresses a

single (but very important) issue: how is electroweak symmetry broken? It does not shed *any* light on the other shortcomings of the SM. For example, SUSY has nothing to say about the pattern of fermion masses and mixings, or the replication of generations. While there are new sources of CP violation in SUSY theories, it is fair to say that SUSY models do not really explain the origin of this. Supersymmetric theories also cause new problems not present in the SM. Why are baryon and lepton number conserved at low energy when it is possible to write dimension four $SU(3) \times SU(2) \times U(1)$ invariant interactions that allow for their non-conservation? Why is the supersymmetric parameter $\mu \sim M_{Weak}$? What is the origin of SUSY breaking and why are SUSY breaking parameters fifteen orders of magnitude smaller than the Planck scale? Why is the CP violation from new SUSY sources so small?

We do not know the answers to these and, probably, several other questions. Perhaps clues to some of these questions lie in the unknown mechanism of SUSY breaking. The measurement of sparticle masses (or other soft SUSY breaking parameters) in future experiments will provide theorists with some guidance in this regard. We should, of course, always keep open the possibility that it is not supersymmetry, but some totally different mechanism that is responsible for stabilizing the electroweak scale. Only experiments can tell whether weak scale supersymmetry is realized in nature. What is clear, however, is that the exploration of the TeV scale will provide essential clues for further unravelling the nature of electroweak symmetry breaking interactions. We must look to see what we find.

Acknowledgements

I thank Dave Soper for his invitation to lecture and participate in the stimulating program of this School. I am grateful to various students for their questions and comments. I also thank K. T. Mahanthappa and his colleagues at Boulder for their generous hospitality as well as for providing a wonderful environment for the School. I am grateful to Howie Baer, Vernon Barger, and especially, Manuel Drees (who also let me know in no uncertain terms what a lousy proof-reader I am) for their comments on this manuscript. I am also grateful to various colleagues for discussions and collaborations on supersymmetry phenomenology. Without their input, these lectures would not have been possible. My thanks also to the members of the Phenomenology Institute, University of Wisconsin, Madison for their wonderful hospitality during the time that these notes were prepared. This research is supported in part by the U.S. Department of Energy grant DE-FG-03-94ER40833.

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